VARIABILITY AND SPECTRAL MODELING OF THE HARD X-RAY EMISSION OF GX 339–4 IN A BRIGHT LOW/HARD STATE

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ABSTRACT

We study the high-energy emission of the Galactic black hole candidate GX 339–4 using INTEGRAL/SPI and simultaneous RXTE/PCA data. By the end of 2007 January, when it reached its peak luminosity in hard X-rays, the source was in a bright hard state. The SPectrometer on INTEGRAL (SPI) data from this period show a good signal-to-noise ratio, allowing a detailed study of the spectral energy distribution up to several hundred keV. As a main result, we report on the detection of a variable hard spectral feature (≥150 keV) which represents a significant excess with respect to the cutoff power-law shape of the spectrum. The SPI data suggest that the intensity of this feature is positively correlated with the 25–50 keV luminosity of the source and the associated variability timescale is shorter than 7 hr. The simultaneous Proportional Counter Array data, however, show no significant change in the spectral shape, indicating that the source is not undergoing a canonical state transition. We analyzed the broadband spectra in the lights of several physical models, assuming different heating mechanisms and properties of the Comptonizing plasma. For the first time, we performed quantitative model fitting with the new versatile Comptonization code BELM, accounting self-consistently for the presence of a magnetic field. We show that a magnetized medium subject to pure non-thermal electron acceleration provides a framework for a physically consistent interpretation of the observed 4–500 keV emission. Moreover, we find that the spectral variability might be triggered by the variations of only one physical parameter, namely the magnetic field strength. Therefore, it appears that the magnetic field is likely to be a key parameter in the production of the Comptonized hard X-ray emission.

Key words: accretion, accretion disks – magnetic fields – methods: observational – radiation mechanisms: general – X-rays: individual (GX 339–4)

Online-only material: color figures

1. INTRODUCTION

GX 339–4 was discovered in the early 1970s by the MIT X-ray detector aboard the OSO-7 mission (Markert et al. 1973). The source is classified as a low-mass X-ray binary (LMXB; Shabbaz et al. 2001) from upper limits on the optical luminosity of the companion star. It is believed to harbor a black hole, for which Hynes et al. (2003) derived a mass function of $5.8 \pm 0.5 M_\odot$. The inclination of the system is yet uncertain. A study of the binary parameters (Zdziarski et al. 2004) revealed a plausible lower limit of $i \geq 45^\circ$, while Cowley et al. (2002) suggested $i \leq 60^\circ$ because the system is not eclipsing. On the other hand, spectral fits of the Fe Kα region with Chandra (Miller et al. 2004b) and XMM-Newton (Miller et al. 2004a; Reis et al. 2008) clearly favor lower values for the inner disk inclination (typically $i \sim 20^\circ$). This may indicate that the inner accretion disk is warped, making it appear at a lower inclination than the orbital plane. Similarly, there is no certain indication of the distance to the source. Resolving the velocity structure along the line of sight, Hynes et al. (2004) obtained a conservative lower limit of $d \geq 6$ kpc. Zdziarski et al. (2004) analyzed the binary parameters of the system and found the most plausible distance to be around 8 kpc. In this work, we use $d = 8$ kpc, $i = 50^\circ$, and, assuming a small companion mass, we infer $M = 13 M_\odot$ from the mass function.

Black hole binaries (BHBs) are powerful engines producing high-energy radiation up to the γ-ray domain. From an observational perspective, they appear in four spectral states, namely, quiescent, low/hard, intermediate, and high/soft (Tanaka & Lewin 1995; Belloni et al. 2005; see also McClintock & Remillard 2006 for a slightly different classification).

Soft states. In the soft state, the energy spectrum is dominated by a soft ($\lesssim 10$ keV) component which is attributed to thermal emission from an optically thick geometrically thin accretion disk (Shakura & Sunyaev 1973) extending down to the last stable orbit. Above 10 keV, the emission is characterized by a complex hard X-ray continuum with no clear presence of a high-energy cutoff (see, e.g., Gierliński et al. 1999; Motta et al. 2009; Caballero-García et al. 2009). This component can be attributed to inverse Compton scattering of soft photons (UV, soft X) in a hybrid thermal/non-thermal electron plasma (the so-called corona; see, e.g., the review by Done et al. 2007, hereafter DGK07).

The non-thermal particles in the corona may be generated due to the magnetic field. Indeed, the Parker instability is able to transport a significant fraction of the accretion power above and below the disk (Galeev et al. 1979; Uzdensky & Goodman 2008), where the energy may then be dissipated through magnetic reconnection. Alternatively, Fermi acceleration at relativistic shocks is also expected to produce non-thermal particle distributions which can explain the observed steep power-law emission in hard X-rays.

Hard states. In the hard state, the energy spectrum is very different. The disk emission almost vanishes and the spectrum is dominated by a hard power-law component (photon index $\Gamma = 1.4–2.0$) with a nearly exponential cutoff ($E_{\text{cut}} = 50–150$ keV; see, e.g., Zdziarski et al. 1998). Such a spectrum is well described by thermal Comptonization in a hot, optically thin electron–proton plasma (Sunyaev & Truemper 1979; Zdziarski & Gierliński 2004). Soft γ-ray observations of several BHBs additionally revealed a high-energy excess with respect to a thermal Comptonization model, suggesting that in some cases
at least some level of non-thermal electron acceleration is also required. For instance, COMPTEL observed a non-thermal tail in the averaged hard state spectrum of Cygnus X-1 (McConnell et al. 2002) while OSSE and SPECTrometer on INTEGRAL (SPI) detected such a feature during bright hard states of GX 339−4 (Johnson et al. 1993; Wardziński et al. 2002; Joinet et al. 2007).

Intermediate states. Intermediate states are observed during transitions between the hard and the soft states and usually show the characteristic features of both (Miyamoto et al. 1991; Mendez & van der Klis 1997; Kong et al. 2002). For GX 339−4, Belloni et al. (2005) defined two different varieties of intermediate state: the hard intermediate state (HIMS) and the soft intermediate state (SIMS). The transition from the HIMS to the SIMS of the 2004 outburst of GX 339−4 is reported in Del Santo et al. (2008).

According to a popular scenario, the different spectral states can be explained through changes in the geometry of the accretion flow. The weakness of the thermal component in the hard state is generally interpreted as a consequence of a truncated accretion disk (DGK07), which, in its inner parts, is replaced by a hot, advection dominated accretion flow (ADAF; Shapiro et al. 1976; Narayan & Yi 1994; Yuan et al. 2007). In these solutions, gravitational energy is converted into thermal energy of protons, which in turn heat the electrons through Coulomb collisions. This process naturally forms the quasi-thermal electron distributions that are required to explain the typical hard state spectra. Moreover, recent models of hot accretion flows include a small non-thermal component that can account for the sometimes observed high-energy excess.

However, a number of recent results seem to question this paradigm. First, it is notoriously difficult to estimate the inner disk radius in the hard state and its recession is a highly debated topic (see, e.g., Cabanac et al. 2009). Although several reports support the truncated disk scenario (Tomsick et al. 2009; Done & Diaz Trigo 2009), others suggest that even in hard states the accretion disk may extend down to a few Schwarzschild radii (Miller et al. 2006; Reis et al. 2008), which would be inconsistent with the standard ADAF scenario. Second, hot accretion flows generally exist only if the optical depth is small ($\tau_T \ll 1$), which implies that the electron–proton coupling is weak and therefore allows high proton temperatures. However, it is difficult to align such a small optical depth with both the hard X-ray spectral slope and the thermal cutoff energy. When it is difficult to align such a small optical depth with both the weak and therefore allows high proton temperatures. However, it is still not straightforward to accommodate the inferred high-energy cutoff to the observations, even if gravitational redshift is accounted for.

As an alternative to the ADAF-like models, the Comptonizing medium in the hard state could as well be powered by the same non-thermal mechanisms that are believed to accelerate the electrons in the soft state, i.e., diffusive shock acceleration or magnetic reconnection. Such a model naturally accounts for the presence of a non-thermal component in the high-energy spectrum. In addition, MB09 showed that the steady state electron distribution can appear quasi-thermal even if acceleration mechanisms are purely non-thermal (see also Poutanen & Vurm 2009). These authors studied the thermalizing effects of the magnetic field, since it was pointed out by Ghisellini et al. (1998) that the very fast emission and absorption of synchrotron photons (the so-called “synchrotron boiler effect”) is able to thermalize the electron distribution in a few light-crossing times. Using BELM, a new code which includes this effect (Belmont et al. 2008), MB09 qualitatively explained the variety of spectral states observed in the prototypical BH binary Cygnus X-1. The model is consistent with a disk recession in the hard state (Poutanen & Vurm 2009), but since the weakness of the soft component could as well result from a lower disk temperature (Beloborodov 1999; Malzac et al. 2001), it does not necessarily require a change in the geometry of the accretion flow (MB09).

In this paper, we use observations of GX 339−4 with INTEGRAL and PCA/RXTE to study the spectral energy distribution of the bright hard state from the 2007 outburst. We analyze the spectral variability in the framework of the models cited here before. The data and the associated reduction procedures are described in Section 2. Section 3 is dedicated to a phenomenological analysis of the high-energy behavior while in Section 4 we present an extensive physical analysis of the broadband spectra. The results and their implications are discussed in Section 5 and summarized in Section 6.

2. OBSERVATIONS

The data presented here on GX 339−4 were obtained with the INTEGRAL and RXTE observatories at the end of 2007 January. At that time, the source was very bright in hard X-rays (cf. Figure 1) and Motta et al. (2009) identified its spectro-temporal properties to be characteristic of the hard state.

2.1. Instruments

Our analysis focuses on the results obtained by the spectrometer SPI (Vedrenne et al. 2003), which is one of the two main instruments aboard the international gamma-ray observatory INTEGRAL. Operating in the 20 keV–8 MeV band, SPI uses germanium detectors to provide high spectral resolution while imaging is performed through the coded mask technique (Roques et al. 2003). The observational strategy of the INTEGRAL mission is to sample the region of interest by means of 30–40 minute long fixed pointings (the so-called science windows), each separated by a 2° angular distance (see Jensen et al. 2003 for details). In order to extend the spectral coverage to lower energies, we also considered simultaneous 4–28 keV data from the Proportional Counter Array (PCA) on RXTE (Bradt et al. 1993).

2.2. Data and Reduction Methods

The SPI data from the GX 339−4 region were obtained during INTEGRAL revolution 525, lasting from 2007 January 30 to February 1 (cf. Table 1). We selected the science windows where GX 339−4 was less than 12° off the central axis and which did not show any contamination by solar flares or radiation belt exit/entry. This resulted in a set of 50 useful pointings representing 107 ks of observational coverage. In order to determine the emitting sources in the field of view, we performed 25–50 keV imaging with the SPIROS software (Skinner & Connell 2003). Aside from the main source, 4U 1700–377, OAO 1657−415, IGR J16318−4848, and GX 340+0 were detected above a 5σ threshold. The positions of these sources were then given as a
Figure 1. Overview of the 2007 outburst showing the daily light curves of GX 339−4 obtained by SWIFT/BAT (15–50 keV; top panel) and RXTE/ASM (1.3–12 keV; lower panel). The vertical dotted lines indicate the time period during which the SPI data presented here were obtained.

(A color version of this figure is available in the online journal.)

Table 1 Log of the Observations Discussed in the Paper

| Instruments       | Obs. ID   | MJD Start | MJD Stop  | Exp Time (ks) |
|-------------------|-----------|-----------|-----------|---------------|
| SPI/IBIS          | 525       | 54130.6   | 54132.2   | 107           |
| SPI/IBIS Low (cutoff) | ...       | ...       | ...       | 36            |
| SPI/IBIS High (excess) | ...       | ...       | ...       | 36            |
| PCA/HEXTE         | 92035-01-01-02 | 54131.10 | 54131.17  | 3.7           |
| PCA/HEXTE         | 92035-01-01-04 | 54132.08 | 54132.15  | 3.7           |

priori information to a specific flux-extraction algorithm, using the SPI instrument response for sky-model fitting. In a first run of the software, we allowed the background normalization as well as the source fluxes to vary between successive pointings. In this way we determined the most appropriate variability timescale for each component. The background normalization was more or less stable during the SPI observation; therefore we assumed no background variability in the final reduction process. Some of the sources in the field, however, showed considerable flux variations. Most notably, 4U 1700−377 was extremely variable, with a 20–50 keV flux per science window ranging from 0 to 700 mcrab.

There are two RXTE observations which are simultaneous with the INTEGRAL exposures (see Table 1). The first observation took place toward the middle and the second toward the end of revolution 525 (cf. the shaded areas in Figure 2). For both observations, we downloaded the standard products from the HEASARC Web site3 and analyzed the data in xspec v11.3.2 (Arnaud 1996). Apart from a global 3% flux difference, the PCA spectra from both observations are compatible within the error bars and were co-added using the addspec routine from the ftools package. Following Motta et al. (2009), we added a systematic error of 0.6% to each PCA channel.

2.3. SPI Light Curve and Data Subsets

Figure 2 shows the 25–50 keV SPI light curve from GX 339−4 obtained during revolution 525. Each bin represents one pointing or science window, which corresponds to a timescale of about 40 minutes. Overall, the 25–50 keV flux shows no particular evolution, indicating that the source remained in the bright hard state during the ∼1.5 day long observation. However, the light curve reveals minor variability (on a timescale of hours, or less) around the average flux value ($\langle F \rangle = 657.1 \pm 3.8$ mcrab). To investigate whether this variability could be linked with changes in the spectral energy distribution, we fixed two arbitrary flux bounds ($F_{\text{low}} = 640$ mcrab and $F_{\text{high}} = 670$ mcrab) and grouped the science windows where $F_{\text{scw}} < F_{\text{low}}$ and $F_{\text{scw}} > F_{\text{high}}$, respectively. For each of the two resulting data subsets, color-coded in blue and red and referred to as low and high respectively, we then produced the averaged 25–500 keV spectrum. Both subsets consist of an equal number of science windows, which allows a straightforward comparison of the spectra. Since the source flux is outstandingly high, the spectral extraction yields significant results, even if the number of science windows is relatively small.

3. HIGH-ENERGY EMISSION

3.1. SPI Spectra

The resulting 25–500 keV spectra differ by ∼14% in terms of their total flux and were separately fitted in xspec with a simple cutoff power law model. The uncertainty on the model parameters is given at the 90% confidence level ($\Delta \chi^2 = 2.7$). As the SPI data alone do not allow us to simultaneously constrain the photon index and the cutoff energy, we fixed the former to $\Gamma = 1.3$ and left the latter free to vary. The model provides a good fit to the low spectrum ($\chi^2/22 = 0.75$) and we infer a significant cutoff at $E_c = 58.6 \pm 2.2$ keV. For the high spectrum, the marked curvature around 40 keV still suggests

3 http://heasarc.gsfc.nasa.gov/docs/archive.html.
the presence of a cutoff energy, fitted at $E_c = 56.2 \pm 2.1$ keV. With respect to the model, however, we observe a significant excess above 150 keV, leading to an unacceptable quality of the fit ($\chi^2/28 = 1.82$). To account for the high-energy excess, we added a second power law of fixed photon index $\Gamma = 1.6$. The presence of this second component, which reduces the cutoff energy to $E_c = 49.5 \pm 3.8$ keV but improves the fit to $\chi^2/27 = 1.03$, is statistically required as shown by the $P_{\text{test}}$ (Bevington & Robinson 1992). Indeed, we infer a probability of $P_{\text{test}} = 6.1 \times 10^{-5}$ that the improvement of the fit was a chance event. As a consequence, we conclude that both SPI spectra essentially differ in terms of the highest energy emission ($>150$ keV).

3.2. Imaging and Spectral Robustness

As the flux extraction process of the SPI data can be sensitive to background modeling and a priori information about the distribution of emitting sources in the field of view, we double checked our results. In the 150–450 keV $\text{spiros}$ image drawn from the high data set, GX 339−4 is detected without any a priori information at a flux level of $347.4 \pm 49.9$ mcrab and a 7.0 $\sigma$ significance. This contrasts to a non-detection in the low data set, where the flux significance at the source position is below 2.5 $\sigma$. We note that the maximum level of the uniformly distributed residuals is around 4$\sigma$ in both images.

Since GX 339−4 is the only source detected above 150 keV, we used a sky model with a single point source and re-extracted the high-energy part ($>150$ keV) of the spectra. We also tested different background models along with various variability timescales. The obtained results are all perfectly compatible within the 1$\sigma$ errors, showing that the spectra are insensitive to the details of the flux extraction process. We conclude that the respective presence of a cutoff and a high-energy excess in the SPI data is robust and significant. As a consequence, we rename the low and high spectra as cutoff and excess, respectively.

3.3. Time Evolution

The presence and absence of a high-energy excess in the high and low flux spectra suggests that the appearance of this feature is correlated with the 25–50 keV flux of the source. To investigate this issue further, we extracted light curves in the 25–50, 50–150, and 150–450 keV energy bands with a binning of 10 science windows ($\approx 7$ hr) and plotted the [150–450]/[50–150] keV hardness ratio as a function of time and as a function of the 25–50 keV flux. From these plots, shown in Figure 3, it is clear that the strength of the high-energy excess is correlated with the 25–50 keV flux. Since a different time binning (8 or 12 science windows for instance) leads to the same results and since the 25–50 keV flux is a good tracer of the total X-ray luminosity, we conclude that the intensity of the high-energy tail is most likely correlated to the total X-ray luminosity of the source.

For statistical reasons, the SPI data do not allow us to quantify the exact timescale of the phenomenon. Nevertheless, from the 25–50 keV science-window light curve (cf. Figure 2) and the above analysis, we can assess that the timescale of the observed spectral variability is shorter than 7 hr.

3.4. Cross-check with Other Instruments

Since the detection of a high-energy tail is a critical issue, we cross-checked the SPI results with other instruments. We analyzed the data obtained simultaneously by the soft gamma-ray imager IBIS/ISGRI (Ubertini et al. 2003) aboard INTEGRAL and the high-energy X-ray timing experiment HEXTE aboard RXTE. The HEXTE spectra from the two simultaneous RXTE observations (cf. Table 1) have been presented by Motta et al. (2009), who kindly provided us the reduced data. Both spectra are compatible within the 1$\sigma$ errors and were co-added to obtain

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Footnote: If this parameter is left free to vary, the fitted 90% confidence interval is $\Gamma \in [0.3, 1.9]$. 
better statistics at high energies. The total SPI and IBIS/ISGRI spectra (averaged over the whole INTEGRAL revolution 525) have been presented by Caballero-García et al. (2009), who report on the detection of a high-energy excess above 150 keV. We re-extracted the IBIS/ISGRI data using the standard OSA 8.0 software package⁵ and jointly fitted the averaged SPI, IBIS/ISGRI, and HEXTE spectra in xspec. We find a good agreement between the three instruments, not only in spectral shape but as well in normalization. In addition, the data from all three instruments confirm the presence of a high-energy excess with respect to a phenomenological cutoff power law and a thermal Comptonization model (cf. Figure 4, Table 2). For the latter, we used eqpair (Coppi 1999, cf. Section 4.2) with non-thermal heating turned off (ln\_th = 0). The reflection amplitude is poorly constrained and was therefore fixed to Ω/2π = 0.35 (average value derived from the physical analysis, cf. Table 3).

The two IBIS/ISGRI spectra averaged separately over the low and high data sets are consistent with the SPI results. However, due to the shorter exposure times, the high-energy excess is no longer required in the ISGRI data. This is not surprising since the sensitivity of ISGRI above 200 keV is lower than the sensitivity of SPI. As a consequence, we did not consider the IBIS/ISGRI data in the remainder of the paper because they do not improve the high-energy constraints with respect to the models. The HEXTE data, on the other hand, are obtained with even shorter exposure times and are thus unsuitable for our group-wise spectral analysis.

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Table 2: Quantification of the High Energy Excess in the Total Averaged SPI, IBIS/ISGRI, and HEXTE Spectra with Respect to a Cutoff Power Law and a Thermal Comptonization Model

| Instrument   | Model    | (χ²/v)_i | (χ²/v)_f | P_{\text{FTEST}} |
|--------------|----------|-----------|-----------|-----------------|
| SPI          | CUTOFFPL | 36/26     | 18/25     | 4 × 10⁻⁵        |
|              | EQPAIR   | 47/26     | 22/25     | 2 × 10⁻²        |
| IBIS/ISGRI   | CUTOFFPL | 32/29     | 26/28     | 2 × 10⁻²        |
|              | EQPAIR   | 39/29     | 25/28     | 5 × 10⁻⁴        |
| HEXTE        | CUTOFFPL | 59/44     | 47/43     | 2 × 10⁻³        |
|              | EQPAIR   | 66/44     | 47/43     | 1 × 10⁻⁴        |
| SPI+IBIS+HEXTE | CUTOFFPL | 145/103   | 105/102   | 1 × 10⁻⁸        |
|              | EQPAIR   | 141/103   | 99/102    | 2 × 10⁻⁹        |

Notes. To account for the high-energy excess, we added a second power-law component with a fixed spectral index of Γ = 2.0 in the phenomenological model, while we allowed for non-thermal heating in the physical model. Each time we give the fit quality before (χ²/v)_i and after (χ²/v)_f adding a dof, along with the F_{\text{FTEST}} probability that the fit improvement was a chance event.

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4. BROADBAND SPECTRAL ANALYSIS

In order to understand the processes that are likely to cause the observed X/γ-ray emission, we analyzed the broadband PCA/SPI spectra in the lights of different physical Comptonization models. Since the spectral variability is only occurring at high energies, each one of the two SPI spectra (cf. Section 3.1) was jointly fitted with the co-added PCA spectrum. A variable multiplicative factor was added to account for cross-calibration uncertainties as well as the luminosity difference between the two SPI spectra. The difference to unity of this factor is never

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⁵ Available at the INTEGRAL Science Data Centre (ISDC).
larger than 8%. Except for slight changes in the $\chi^2$ values, neither the qualitative nor the quantitative results are affected when using one of the single PCA spectra instead of the co-added spectrum.

4.1. General Model

The high-energy source in GX 339−4 is modeled by a spherical, magnetized, fully ionized proton–electron/positron plasma of radius $R$. The emission of the plasma is derived by self-consistent computations of the equilibrium electron distribution accounting for Compton scattering, synchrotron emission/absorption, pair production/annihilation, and Coulomb collisions ($e-e$ and $e-p$). The Thomson optical depth of the plasma is given by $\tau_T = \tau_{ion} + \tau_{pair}$, where $\tau_{ion}$ represents the optical depth of ionization electrons and $\tau_{pair}$ is the opacity that arises from pair production.

The plasma properties essentially depend on the magnetic field strength as well as on the power supplied by external sources. The tangled magnetic field strength is parameterized by the magnetic compactness:

$$l_B = \frac{\sigma_T}{m_e c^2} R \frac{B^2}{8\pi}$$

(1)

Figure 4. Total averaged SPI (red), IBIS (blue), and HEXTE (green) spectra, jointly fitted with a thermal (left) and a hybrid thermal/non-thermal (right) Comptonization model. The normalization was fixed to the calibration of SPI and we applied multiplicative constants of $C_{IBIS} = 1.03$ and $C_{HEXTE} = 1.10$ to the IBIS/ISGRI and HEXTE spectra, respectively. The lower panels indicate the data/model ratios for each instrument separately. These show that a pure thermal Comptonization model is not able to account for the emission above 150 keV.

Table 3

| Model  | Spec | $\Gamma_{inj}$ | $\gamma_{max}$ | $l_{nth}/l_{B}$ | $l_B$ | $\tau_{ion}$ | $\tau_T$ | $K_{Te}$ (keV) | $\Omega/2\pi$ | $\xi$ | $\chi^2$ (dof) |
|--------|------|----------------|----------------|------------------|-------|--------------|----------|-----------------|-------------|------|----------------|
| ECM1 Cutoff | 3.95$^{+1.55}_{-0.70}$ | 1000* | 4.35 ± 0.07 | $\ldots$ | 2.98 ± 0.07 | 2.99 ± 0.07 | 20.1 | 0.26$^{+0.03}_{-0.03}$ | 710 ± 120 | 0.84(68) |
| Excess | 2.55$^{+0.35}_{-0.15}$ | 1000* | 5.06$^{+0.10}_{-0.08}$ | $\ldots$ | 3.31$^{+0.07}_{-0.15}$ | 3.87$^{+0.07}_{-0.15}$ | 14.3 | 0.25$^{+0.02}_{-0.03}$ | 600 ± 240 | 1.10(74) |
| ECM2 | 2.5* | 4.0$^{+1.3}_{-0.3}$ | 4.25$^{+0.08}_{-0.10}$ | $\ldots$ | 3.06$^{+0.09}_{-0.06}$ | 3.07$^{+0.09}_{-0.06}$ | 19.1 | 0.23$^{+0.02}_{-0.03}$ | 850$^{+140}_{-200}$ | 0.85(68) |
| Excess | 2.5* | 21.6$^{+39}_{-39}$ | 4.94$^{+0.11}_{-0.10}$ | $\ldots$ | 3.69$^{+0.04}_{-0.13}$ | 3.87$^{+0.04}_{-0.13}$ | 14.5 | 0.21$^{+0.03}_{-0.02}$ | 1240$^{+290}_{-240}$ | 1.00(74) |
| ICM1 Cutoff | 3.47$^{+0.36}_{-0.31}$ | 1000* | $\ldots$ | 740$^{+200}_{-140}$ | 2.26 ± 0.18 | 2.26 ± 0.18 | 34.2 | 0.43 ± 0.05 | 300$^{+120}_{-100}$ | 0.96(68) |
| Excess | 2.50$^{+0.32}_{-0.08}$ | 1000* | 25.2 ± 3.6 | 2.14$^{+0.14}_{-0.17}$ | 2.45$^{+0.14}_{-0.17}$ | 25.7 | 0.44$^{+0.06}_{-0.03}$ | 260$^{+110}_{-80}$ | 1.06(74) |
| ICM2 | 2.5* | 15.2$^{+5.5}_{-3.3}$ | $\ldots$ | 283$^{+43}_{-36}$ | 2.30$^{+0.17}_{-0.16}$ | 2.30$^{+0.17}_{-0.16}$ | 33.4 | 0.43$^{+0.04}_{-0.03}$ | 300$^{+120}_{-100}$ | 0.96(68) |
| Excess | 2.5* | 190$^{+110}_{-105}$ | $\ldots$ | 27.3$^{+40}_{-34}$ | 2.49 ± 0.18 | 2.73 ± 0.18 | 23.9 | 0.39$^{+0.04}_{-0.04}$ | 310$^{+155}_{-126}$ | 0.92(74) |

Notes. Fixed parameters are indicated by a star (*) next to their value. The temperature of the thermal component is estimated at the equilibrium using Equation (2.8) of Coppi (1992) and is hence not a fit parameter.
Energy injection is quantified by the compactness parameter:

\[ l = \frac{\sigma_T}{m_e c^3} \frac{L}{R}, \tag{2} \]

where \( L \) is the power supplied to the plasma, \( \sigma_T \) is the Thomson cross section, and \( m_e \) is the electron rest mass. The general model comprises three possible channels for providing energy to the coupled electron–photon system: (1) non-thermal electron acceleration \( \left( l_{\text{nth}} \right) \), (2) thermal heating of the electron distribution \( \left( l_{\text{th}} \right) \), and (3) external soft radiation from a geometrically thin accretion disk \( (l) \). The non-thermal acceleration processes are mimicked by continuous electron injection at a power-law rate with index \( \Gamma_{\text{inj}} \) (i.e., \( n(\gamma) \propto \gamma^{-\Gamma_{\text{inj}}} \)) and Lorentz factors ranging from \( \gamma_{\text{min}} \) to \( \gamma_{\text{max}} \). The thermal heating of the electron distribution could be caused by Coulomb interactions with a population of hot ions. The incident radiation from an accretion disk is modeled by a blackbody component of temperature \( kT_{\text{bb}} \). Since the temperature of this component cannot be constrained by our data, we follow Del Santo et al. (2008) and fix \( kT_{\text{bb}} \) to the fiducial value of 300 eV. All the injected energy ends up in radiation, and in the compactness formalism the total radiated power in steady state holds: \( l = l_{\text{nth}} + l_{\text{th}} + l_\ell \).

For a given source flux \( F \), the total compactness can be estimated by the formula:

\[ l = 100 \left( \frac{F}{F_0} \right) \left( \frac{d}{8 \text{ kpc}} \right)^2 \left( \frac{13 M_\odot}{M} \right) \left( \frac{30 R_G}{R} \right), \tag{3} \]

where \( M \) is the mass of the black hole and \( d \) is the distance to the source. Here we expressed \( l \) in terms of \( F_0 = 2.6 \times 10^{-8} \) erg s\(^{-1}\) cm\(^{-2}\), the observed average 4–500 keV flux of GX 339–4. Assuming a spherical plasma of radius \( R = 30 R_G \), a black hole mass of \( M = 13 M_\odot \), and a distance of \( d = 8 \) kpc, this leads to a total compactness of \( l \sim 100 \).

Besides being a source of soft seed photons, the accretion disk may also give rise to Compton reflection of the incident hard X-rays from the Comptonizing plasma. The reflected emission is calculated using the viewing-angle-dependent Green’s functions (Magdziarz & Zdziarski 1995), but accounting for general relativistic effects. The spatial extension of the possibly ionized disk is parameterized by its inner and outer radii, fixed at \( R_a = 6 R_G \) and \( R_{\text{out}} = 400 R_G \).

Moreover, the model takes into account the K\(_\alpha\) fluorescence of the Fe elements in the disk (diskline; Fabian et al. 1989) as well as the interstellar absorption (phabs). For the latter, we follow Reis et al. (2008) and assume a fixed neutral hydrogen column density of \( N_{\text{H}} = 5.2 \times 10^{21} \) cm\(^{-2}\).

Since our data sets are energy and sensitivity limited and since such a complex modeling is partly degenerate, spectral fitting is unable to provide simultaneous constraints to all the parameters. Therefore, starting from the general model, we define several canonical sub-models where some of the parameters are held invariant.

4.2. Thermal Heating

Given that the relative ratio of the heating mechanisms (thermal versus non-thermal) is very difficult to constrain, we investigated the extreme situations in which there is only one possible channel for providing energy to the electrons.

First, we consider the case where the energy injection is purely thermal \( \left( l_{\text{nth}} \neq 0, l_{\text{th}} = 0 \right) \). The seed photons may arise either from an external source (i.e., the accretion disk), or, in the presence of a magnetic field, from internally generated synchrotron emission. However, since the magnetic field is believed to give rise to non-thermal electron acceleration, we limited our analysis to the standard case where magnetic processes are neglected \( \left( l_{\ell} = 0 \right) \).

The main model parameters are the power supplied to the Comptonizing electrons \( l_{\ell} \), the power contained in the soft photons from the disk \( l_\ell \), the Thomson optical depth of the plasma \( \tau_{\text{th}} \), the reflection amplitude \( \Omega/2\pi \), the ionization parameter of the reflecting material \( \xi \), and the normalization factor between the two instruments. To calculate the model spectra, we use the hybrid thermal/non-thermal Comptonization code EPAIR (Coppi 1999). A detailed description of the code and its application to Cygnus X-1 can be found in Gierliński et al. (1999).

Since the model spectrum strongly depends on the ratio \( l_{\text{th}}/l_\ell \), but only weakly on the absolute values of the compactness parameters, we fix \( l_\ell \) to a constant value in order to improve the fitting strategy. For \( l_\ell \) in the range \( 0.5 < l_\ell \leq 100 \), the spectral shape allows us to constrain \( l_{\text{th}}/l_\ell \) to be always of the order of 5. Therefore we set \( l_\ell = 15 \), so that the total compactness parameter \( l = l_{\text{th}} + l_\ell \) is consistent with Equation (3).

For the cutoff spectrum, however, the thermal model is not appropriate. The best fit is indeed of poor quality, with a reduced chi-square of \( \chi^2/\nu = 1.40 \). We note that the degradation results only from the high-energy channels, which are responsible for \( \chi^2/\nu = 65/23 \) of the total \( \chi^2 \). Obviously, the model fails to reproduce the high-energy excess, supporting the fact that the emission above 100 keV is linked to non-thermal processes. We conclude that a pure thermal model is not adequate to explain the observed spectral behavior.

4.3. Non-thermal Acceleration with External Comptonization

Now we investigate the opposite case, in other words we assume that the energy injection is purely non-thermal \( \left( l_{\text{th}} \neq 0, l_{\text{nth}} = 0 \right) \). The rest of the model remains unchanged, i.e., we consider an external source of soft photons (of fixed compactness \( l_\ell = 15 \)) and we neglect the magnetic field \( \left( l_{\ell} = 0 \right) \). Since the electrons can partly thermalize through Coulomb collisions, the equilibrium distribution is hybrid (i.e., thermal/non-thermal), even if the whole power is supplied via non-thermal injection. The acceleration processes are phenomenologically described by the model parameters \( l_{\text{inj}}, \Gamma_{\text{inj}}, \gamma_{\text{min}}, \text{ and } \gamma_{\text{max}} \). Here again, we define two sub-models in order to reduce the number of free parameters. In the first model, we set \( \left( \gamma_{\text{min}}, \gamma_{\text{max}} \right) = (1.3, 1000) \), and analyze the spectra in terms of \( l_{\text{inj}} \) and \( \Gamma_{\text{inj}} \). This configuration is frequently used in the literature and will provide a familiar context to discuss our results. In the second model, we adopt a more novel approach and investigate the effects of a varying maximum energy of the accelerated electrons. We keep \( \gamma_{\text{min}} = 1.3 \) but fix the spectral index to the fiducial value \( \Gamma_{\text{inj}} = 2.5 \), while \( l_{\text{inj}} \) and \( \gamma_{\text{max}} \) are free to vary. For convenience, we call these models the ECM1 and the ECM2 (for External Comptonization Models), respectively. The associated model spectra are computed with the Comptonization code EPAIR (Coppi 1999) and the fitting results are summarized in Table 3.
4.3.1. The Injection Index

For the cutoff spectrum, the ECM1 provides a good fit to the data ($\chi^2/68 = 0.84$, cf. Figure 5, left). The observed cutoff shape is well reproduced by a very soft injected electron distribution. Our best fit yields a spectral index of $\Gamma_{\text{cutoff}} \sim 0.7$, which was fixed to this value to determine the error on the other free parameters.

We find a compactness ratio of $l_{\text{nth}}/l_e = 4.35 \pm 0.08$, implying a non-thermal compactness of $l_{\text{nth}} = 65.0 \pm 1.0$. The optical depth of ionization electrons is fitted at $\tau_{\text{ion}} = 2.97 \pm 0.07$. Because of the soft injected electron distribution, there is only very little pair production, increasing the total optical depth to $\tau_T \sim 2.99 \pm 0.07$. The observed spectrum strongly requires a moderate amount of Compton reflection, with a fitted amplitude of $\Omega/2\pi = 0.26^{+0.03}_{-0.04}$ and an ionization factor of $\xi = 710 \pm 120$. Freezing $\Omega/2\pi$ to zero leads to a dramatically worse fit ($F$-test probability $< 10^{-20}$).

For the excess spectrum, the best fit ($\chi^2/74 = 1.10$, cf. Figure 5, right) requires a substantially harder power-law injection, with a fitted index of $\Gamma_{\text{inj}} = 2.55^{+0.35}_{-0.15}$. With respect to the cutoff spectrum, the compactness ratio increased by $\sim 15\%$ to $l_{\text{nth}}/l_e = 5.06^{+0.10}_{-0.08}$. The non-thermal compactness now yields $l_{\text{nth}} = 75.9^{+1.5}_{-1.2}$ and the total optical depth is increased to $\tau_T = 3.87^{+0.07}_{-0.05}$, which is mainly due to the enhanced production of pairs. The fitted reflection amplitude and disk ionization, however, remain stable ($\Omega/2\pi = 0.27^{+0.02}_{-0.03}$, $\xi = 600 \pm 240$).

In order to explore the dependence on compactness, we abandoned our fiducial hypothesis $l_e = 15$ and fitted the excess spectrum with a variable soft compactness. We find 90% confidence intervals of $0.3 < l_e < 250$ and $4.5 < l_{\text{nth}}/l_e < 5.9$, which confirms that the compactness ratio does not depend on the total compactness of the source. Although our analysis is unable to uniquely determine the total compactness, the upper and lower limits on the allowed range are well constrained by the spectral shape (Gierliński et al. 1999). The upper bound represents the limit at which, despite an extremely soft injected electron distribution ($\Gamma_{\text{inj}} = 4.50^{+0.03}_{-0.04}$), the plasma is completely pair dominated ($\tau_{\text{ion}} \sim 10^{-4}$; $\tau_T \sim 3.0$; $kT_e \sim 11.5$ keV). Above this threshold, the growing amount of pairs can no longer be balanced by a softer injection spectrum, leading to an inaccurate reproduction of the thermal peak in the photon spectrum (cf. Section 5.2). The lower bound corresponds to the limit at which the cooling of the high-energy particles is dominated by Coulomb instead of Compton losses, leading to an underprediction of the intensity of the high-energy tail. A change from $l_e = 0.3$ to $l_e = 0.25$ implies a degradation of $\Delta \chi^2 = 5$ at 74 degrees of freedom (dof), showing that the threshold is well established. Hence, in the framework of an external Comptonization model involving non-thermal electron acceleration up to $\gamma_{\text{max}} = 1000$, the total compactness of the hard X-ray source can be conservatively constrained by means of pure spectral analysis to be $2 < l < 1500$.

4.3.2. The Maximum Particle Energy

Now, in the ECM2, we freeze $\Gamma_{\text{inj}}$ to 2.5 but allow for variations of the maximum Lorentz factor of the accelerated electrons. For the cutoff spectrum, we obtain again a very good fit to the data ($\chi^2/68 = 0.85$). The effects of a much harder slope are effectively balanced by a very low cutoff energy of the injected electron distribution. Indeed, we find $\gamma_{\text{max}} = 4.0^{+0.3}_{-0.2}$, which means that the maximum difference of the kinetic energy between the accelerated particles is about a factor of 10. The other parameters are not much affected, namely we find $l_{\text{nth}}/l_e = 4.25^{+0.08}_{-0.10}$ and $\tau_{\text{ion}} \sim \tau_T = 3.06^{+0.09}_{-0.06}$.

Figure 5. Joint SPI and PCA spectra fitted with the ECM1. The total model is shown by the solid line, the dot-dashed line marks the iron fluorescence emission from the disk, the dashed line shows the reflection component, and the dotted line represents the sum of the Comptonized and reflected emission. The error bars are given at the 1σ level.

(A color version of this figure is available in the online journal.)
the reflection parameters remain stable ($\Omega/2\pi = 0.23^{+0.03}_{-0.16}$, $\xi = 850^{+140}_{-200}$).

For the excess spectrum, the fixed injection index is consistent with the best-fit value obtained with the ECM1. Nevertheless, it turns out that there are some differences between the results obtained with the two models. Mainly, we note that particle acceleration up to $\gamma_{\text{max}} = 21.6^{+39.0}_{-7.0}$ is sufficient to reproduce the observed high-energy tail. Moreover, a truncated electron distribution allows us to improve the quality of the best fit to $\chi^2/\nu = 1.00$. The total opacity of the plasma is found equal to the value obtained with a non-truncated distribution, but the pair yield is reduced due to the lack of very energetic particles. With respect to the ECM1, the other parameters are only marginally affected and remain consistent within the 90% confidence errors.

### 4.4. Non-thermal Acceleration with Internal Comptonization

Finally, we consider models in which the observed X/$\gamma$-ray spectra are produced in a magnetized plasma. To emphasize the effects of the magnetic field, we study the case where all the seed photons are internally generated from synchrotron emission ($l_s = 0$, $l_B \neq 0$). As in the previously analyzed non-magnetic models, we adopt the same configurations to phenomenologically describe the acceleration mechanism. For convenience, we call these models the ICM1 and the ICM2 (for internal Comptonization models), respectively. To calculate the model spectra, we used the new versatile Comptonization code \texttt{BELM} (Belmont et al. 2008). One of the main differences with \texttt{EPAIR} is that \texttt{BELM} accurately accounts for self-absorbed cyclo-synchrotron radiation from the sub-relativistic to the ultra-relativistic regime. We compared the spectra obtained by both codes for the best-fit parameters of the non-magnetized models. The relative differences are smaller than 3% at all energy. The fitting results are summarized in Table 3.

#### 4.4.1. The Injection Index

First, we employ the commonly used configuration, i.e., we fix ($\gamma_{\text{min}}, \gamma_{\text{max}}) = (1.3, 1000)$. A qualitative analysis shows that the model spectrum below 30 keV is rather insensitive to the individual values of $l_{\text{nth}}, l_B$, and $\Gamma_{\text{inj}}$, but strongly depends on a combination of all three parameters (see also MB09). The situation is similar as in the non-magnetic case (cf. Section 4.2), although the dependence is slightly more complicated. On the other hand, the high-energy spectrum (>100 keV) is mostly determined by the injection index $\Gamma_{\text{inj}}$. As a consequence, we use Equation (3) to fix $l_{\text{nth}} = 100$, while the broadband spectra allow us to disentangle the degeneracy between $l_B$ and $\Gamma_{\text{inj}}$.

Due to the detailed treatment of the microphysics in the \texttt{BELM} code, real-time fits in \textsc{xspect} are very time consuming. In order to make the fitting process more efficient, we tabulated the model spectra.\footnote{We used the \texttt{WFTBMD} routine (publicly available at the HEASARC website) to create the appropriate FITS file required by the \texttt{ATABLE} model in \textsc{xspect}.} The resulting table file has three dimensions (corresponding to the parameters $l_B, \Gamma_{\text{inj}}$, and $\tau_{\text{ion}}$) and the fits are performed through interpolation between the tabulated spectra.

In order to account for Compton reflection from a cold disk, the \texttt{BELM} model is convolved with an angle-dependent reflection routine based on the \texttt{PEKIRV} model by Magdziarz & Zdziarski (1995), but taking into account the distortions due to general relativistic effects. The free parameters of the ICM1 are hence $l_B, \Gamma_{\text{inj}}, \tau_{\text{ion}}, \Omega/2\pi, \xi$, the iron line energy, and the normalization factor between PCA and SPI.

For the cutoff spectrum, the ICM1 provides a good fit to the data ($\chi^2/\nu = 68 = 0.96$). In comparison with the ECM1, the high-energy rollover is better reproduced. However, as can be seen from the residuals in Figure 6 (left), the fit is slightly less accurate in the iron line region around 6.4 keV. There are also some positive residuals below 5 keV, possibly hinting at the need for a soft disk component. Since we focus on the...
spectral behavior at high energies, we did not investigate these issues any further. The best fit is obtained with an electron spectral index of $\Gamma_{\text{inj}} = 3.47^{+0.66}_{-0.31}$. The magnetic compactness is fitted at $l_B = 740^{+200}_{-140}$, which corresponds to a magnetic field strength of $B = 2.0^{+1.0}_{-0.9} \times 10^7$ G. The plasma is found to be of moderate optical thickness, with fitted $\tau_{\text{ion}} = 2.26 \pm 0.18$. Due to the fast synchrotron cooling of the small number of high-energy particles, the produced pair yield is negligible. As with the ECM1, the data strongly require ionized Compton reflection, with a fitted amplitude and ionization factor of $\Omega/2\pi = 0.43 \pm 0.05$ and $\xi = 300^{+120}_{-100}$, respectively.

For the excess spectrum, the ECM1 again allows a good description of the data (cf. Figure 6, right). The best fit yields $\chi^2/\nu = 1.06$ and the high-energy data constrain the injection index to $\Gamma_{\text{inj}} = 2.50^{+0.38}_{-0.08}$. This result is roughly equal to the value obtained with the ECM1. However, at equal injected electron distribution, the high-energy tail of the ICM1 spectrum is slightly steeper. Indeed, contrary to a non-magnetized model, the Compton losses compete with the synchrotron losses and a significant fraction of the energy radiated by the non-thermal leptons is emitted in the optical/UV range rather than in hard X-rays. Accounting for the high-energy tail therefore requires more high-energy particles, i.e., a smaller $\Gamma_{\text{inj}}$. We find a magnetic compactness of $l_B = 25.2 \pm 3.6$, which given our assumptions translates to $B = 3.76 \pm 0.14 \times 10^6$ G. This means that in the framework of this model, the magnetic field strength would drop by a factor of 5.5 during a change from the cutoff to the excess spectrum.

The Thomson optical depth of ionization electrons is fitted at $\tau_{\text{ion}} = 2.14^{+0.14}_{-0.17}$. Contrary to the results obtained with the ECM1, the $\tau_{\text{ion}}$ values for both spectra are compatible within the 90% confidence errors. The reflection routine yields $\Omega/2\pi = 0.44^{+0.06}_{-0.03}$ and $\xi = 260^{+110}_{-80}$, which means that these characteristics did not change either.

4.4.2. The Maximum Particle Energy

In the ECM2, we keep $l_{\text{th}} = 100$ and $\gamma_{\text{min}} = 1.3$, but allow now for variations of $\gamma_{\text{max}}$. Instead, we set the injection index to $\Gamma_{\text{inj}} = 2.5$ and generate a new fitting table which again has three dimensions, corresponding to the free parameters $l_B$, $\gamma_{\text{max}}$, and $\tau_{\text{ion}}$.

For the cutoff spectrum, we obtain a good fit to the data ($\chi^2/\nu = 0.96$; shown in Figure 7, left), qualitatively equivalent to the best fit obtained with the ECM1. Again, this shows that the effects of a very soft injection slope can be mimicked by a much harder injected distribution, but truncated at a certain particle energy. We find $\gamma_{\text{max}} = 15.2^{+5.5}_{-3.0}$, which is significantly higher than the fitted maximum particle energy in the ECM2. The inferred magnetic compactness is reduced with respect to the ECM1, that is to say $l_B = 280^{+43}_{-36}$, which for a medium of typical size $R = 30 R_G$ corresponds to $B = 1.25^{+0.5}_{-0.4} \times 10^7$ G. The other parameters are not much affected, namely we find $\tau_{\text{ion}} \simeq 30.3^{+0.16}_{-0.16}$ and from the reflected component we infer $\Omega/2\pi = 0.43^{+0.04}_{-0.03}$ and $\xi = 300^{+120}_{-100}$.

For the excess spectrum, we note that the injection index is fixed to the best-fit value obtained with the ECM1. However, allowing for $\gamma_{\text{max}}$ to vary (i.e., allowing for a truncated electron distribution), the fit may be improved considerably. For $\gamma_{\text{max}} = 190^{+110}_{-105}$, we obtain the best description of the excess spectrum, with a reduced $\chi^2$ of $\chi^2/\nu = 0.92$ (cf. Figure 7, right). Compared to the non-magnetized model, the maximum electron energy is again found to be much higher. This is expected since in the magnetic models, the seed photons have a lower average energy than in the non-magnetized models. Indeed, in the ICMs, the high-energy photons are produced from single Compton scattering off the synchrotron emission ($E_s \sim 0.01$ keV) while in ECMs the seed photons originate from the disk emission ($E_s \sim 1$ keV). Therefore, in order to upscatter the seed photons

Figure 7. Joint SPI and PCA spectra fitted with the ICM2. See Figure 6 for the description of the different model components. The error bars are given at the 1σ level.
(A color version of this figure is available in the online journal.)
to 200 keV, the electrons must have averaged Lorentz factors of $\gamma \sim 12$ in the ECMs and $\gamma \sim 120$ in the ICMs.

The magnetic compactness remains equal to the value inferred with the ECM1, namely $l_B = 27.3_{-4.3}^{+4.0}$. A transition from the cutoff to the excess spectrum thus requires a factor 3.2 decrease of the magnetic field strength. The opacity from ionization electrons yields $\tau_{\text{ion}} = 2.49 \pm 0.18$, while the total optical depth is found to be $\tau_T = 2.73 \pm 0.18$. This shows again that the ionization opacity does not change between the two spectra. Finally, the fitted reflection amplitude and ionization parameter are found to be consistent for both spectra regardless of the acceleration model (cf. Table 3). We thus conclude that in the framework of a strongly magnetized medium, the constraints to these parameters are very robust.

5. DISCUSSION

The SPI observations showed that during the bright hard state of the 2007 outburst, the highest energy emission (>150 keV) of GX 339–4 was variable. While the spectral shape at lower energies (4–150 keV) remained more or less constant, we detected the significant appearance/disappearance of a high-energy tail. The strength of this hard tail, varying on a timescale of less than 7 hr, is found to be positively correlated with the total X-ray luminosity of the source and enables interesting constraints to the physical processes which could be responsible for the high-energy emission.

5.1. The Pure Thermal Model

The clear detection of a cutoff energy in the cutoff spectrum indicates that the Comptonizing electron distribution is quasi-Maxwellian. Thus, it is not surprising that the spectrum can be explained by assuming only thermal heating of the plasma. As suggested by the ADAF models, this could be achieved through Coulomb interactions with a thermal distribution of hot protons. The temperature of the protons can be estimated from the thermal compactness, the electron temperature, and the optical depth of the plasma (cf. formula (4) in MB09). For $l_{\text{th}} = 100$, we infer a proton temperature of the order of 1 MeV. This is of the same order of magnitude as the proton temperature estimated in MB09 for the canonical hard state of Cygnus X-1, which in comparison to the hard state of GX 339–4 analyzed here shows a hotter electron plasma ($kT_e \sim 85$ keV) but a lower compactness ($l_{\text{th}} \sim 5$). As mentioned in MB09, proton temperatures of the order of 1 MeV are significantly lower than what is expected in typical two-temperature accretion flows, namely $kT_i > 10$ MeV.

In any case, pure thermal heating is not enough since it is not able to explain the appearance of the observed high-energy tail. We conclude that either the hard excess is independent of the thermal component, in which case its origin is located outside the innermost regions, or that both components are linked, in which case at least some level of non-thermal heating is required.

On the other hand, both broadband spectra can be successfully explained by models involving only non-thermal electron acceleration. This suggests that the Comptonizing medium in hard states could be powered by the same non-thermal mechanisms that are believed to power the accretion disk corona in soft states. Depending on the nature of the plasma (magnetized or not) and the origin of the seed photons, these issues are discussed in the next paragraphs.

5.2. The Non-magnetic Case

In the framework of a non-magnetized model, the variability of the high-energy spectrum can be explained by changes in the properties of the involved acceleration processes. Namely, the fits suggest a small variation of the total power supplied to the plasma along with a significant variation of either the spectral index (in the ECM1) or the cutoff energy (in the ECM2) of the non-thermal electron distribution.

In the ECM1, a change from the cutoff to the excess spectrum requires that the spectral index typically drops from $\Gamma_{\text{inj}} = 4.0$ to 2.5. This implies that the average energy of the accelerated particles ($\langle E_{\text{inj}} \rangle$) rises from 1.0 to 1.8 MeV. In the ECM2, at constant spectral index, the best fits show that relatively low maximum Lorentz factors are sufficient to reproduce the data. In this case, the appearance of the high-energy excess requires an increase of $\gamma_{\text{max}}$ from 4 to 22, which implies that the average energy of the accelerated electrons rises accordingly, from 1.0 to 1.5 MeV. In both cases, such variations can be explained by the possible non-stationarity of the inherent acceleration mechanisms. For instance, in the framework of shock acceleration, the properties of the accelerated particles depend on the shock strength (Webb et al. 1984; Spitkovsky 2008). Acceleration by reconnection depends on several physical parameters such as the local geometry of the reconnection zone (Zenitani & Hoshino 2007) and the number of reconnection sites (if the particles are accelerated stochastically by successive acceleration events in different sites (Anastasiadis et al. 1997; Dauphin et al. 2007)). All these properties may undergo variations with time and could hence explain the observed variability.

Our results are consistent with those of Gierliński & Zdziarski (2005), who studied the energy-dependent variability of non-magnetized Comptonization models in response to varying physical parameters. Although the variations of $\Gamma_{\text{inj}}$ and $\gamma_{\text{max}}$ could not explain the observed rms spectra of XTE J1650–500 and XTE J1550–564, they found that changes in these parameters produce strong variations in the X-ray spectrum above 50 keV, which is what is required here to reproduce the present data.

The spectral analysis suggests that a transition from the cutoff to the excess spectrum additionally requires a 15% increase of the total power supplied to the plasma. Considering a spherical medium of fixed radius and constant illumination from the accretion disk, this increase is consistent with the 14% difference in the observed 4–500 keV luminosity. In addition, the total optical depth of the plasma is found to increase by about 30%, which is mainly due to the enhanced pair production occurring in the harder acceleration regime.

Motivated by size and luminosity estimates, we assumed a constant soft photon compactness of $l_s = 15$. As mentioned earlier, the above results are only weakly dependent on the exact value of $l_s$. Very high or very low values of the illumination compactness, however, turned out to be inconsistent with the observed spectra. Indeed, since the ratio $l_{\text{nth}}/l_s$ is robustly constrained, strong illumination requires an efficient acceleration mechanism (i.e., large $l_{\text{nth}}$) to reproduce the spectral slope at lower energies (< 20 keV). Consequently, this generates very energetic radiation which produces large amounts of $e^{-} e^{+}$ pairs through photon–photon annihilation. This, in turn, reduces the equilibrium temperature of the plasma since more particles have to share the same amount of energy. Hence, above a certain soft photon compactness, the equilibrium temperature will be too low to be consistent with the observed thermal peak of the spectrum.
Reciprocally, since the thermal part of the electron distribution is roughly determined by the balance between acceleration and Compton cooling, decreasing $l_\gamma$ at constant $l_{nth}/l_\gamma$ has no significant effect on the lower energy part of the photon spectrum. However, a major fraction of the high-energy tail results from the Comptonization by mildly relativistic particles ($2 < \gamma < 10$). Below a certain soft photon compactness, the cooling of these mildly relativistic particles is no longer dominated by the Compton losses but by Coulomb interactions with the lower-energy thermal electrons. Thus, decreasing $l_\gamma$ at constant $l_{nth}/l_\gamma$ provides a weaker acceleration rate while the cooling remains constant. This reduces the intensity of the high-energy tail up to the point where the model predictions are no longer consistent with the data.

In conclusion, independently of any geometric argument, we obtain conservative bounds to the total compactness of the X-ray-emitting plasma, i.e., $2 < l < 1500$. These limits are consistent with the estimates derived from geometric arguments, but unfortunately not very constraining. Anyhow, the robustness of the fitted compactness ratio $l_{nth}/l_\gamma \simeq 4.5$ allows us to conclude that the luminosity of the cold disk represents at most $\sim 20\%$ of the luminosity of the Comptonized component, possibly much less if the plasma is magnetized.

5.3. The Magnetic Case

Using the new code belm (Belmont et al. 2008), we showed that the hard X-ray behavior of GX 339–4 in the bright hard state can be explained by assuming pure non-thermal electron acceleration and subsequent Comptonization of the self-consistently produced synchrotron photons. The model requires no incident radiation from the accretion disk and assumes constant power injection into the magnetized plasma. As in the non-magnetic case, the spectral variability can be mimicked by two different configurations of the acceleration model, involving either a variable power-law slope (in the ICM1) or a variable maximum energy (in the ICM2) of the injected electron distribution.

In principle, these models allow us to estimate the averaged magnetic field strength of the plasma. However, since the fits provide precise constraints only for the ratio $l_B/l\gamma$, the uncertainties on the total compactness (cf. Equation (3) are projected to the estimate of $l_B$. To discuss our results, we express $l_B$ as a fraction of the magnetic compactness at equipartition with the radiation field $l_{B_0}$. As we have the approximate dependence $l_{B_0} \propto l \times (1 + \tau_f/3)$ (cf. Equation (8) in MB09), the ratio $l_B/l_{B_0}$ does not depend on the uncertainties regarding the source size and distance and is therefore a good indicator to quantify the importance of the magnetic processes. In addition, $l_\gamma = 0$ was assumed to study the physics of a strongly magnetized medium, but it cannot be excluded that both the synchrotron flux and a soft disk component contribute to the cooling of the non-thermal particles. If the medium is additionally illuminated by cold disk photons, less synchrotron cooling will be required to reproduce the slope of the lower energy spectrum, implying that the fitted values of the magnetic compactness are in fact conservative upper limits. In the ICM1, we infer $l_B/l_{B_0} < 18$ from the cutoff and $l_B/l_{B_0} \leq 0.58$ from the excess spectrum. In the ICM2, the fitted magnetic compactness for the cutoff spectrum is lower, namely we find $l_B/l_{B_0} \leq 6.0$ and $l_B/l_{B_0} \leq 0.60$, respectively.

As mentioned above, the magnetic models do not require any disk blackbody photons to produce the observed 4–500 keV spectra. If the accretion disk extends down close to the black hole (as requested by the accretion disk corona models), the paucity of soft disk photons can be explained if the Comptonizing coronal material is outflowing at a mildly relativistic speed (Beloborodov 1999; Malzac et al. 2001). Using the formulae (5) and (7) of Beloborodov (1999), we estimate that bulk velocities of at least $0.6 c$ are required for the cooling of the corona being dominated by synchrotron self-Compton. Comptonization off a dynamical corona may then blueshift the emerging spectrum, but these corrections remain moderate at $\sim 0.6 c$ and have not been included in the spectral fits. On the other hand, the disk may as well be truncated (as requested by the hot flow models), since the data do not explicitly require relativistic smearing of the reflection features. Thus, as long as the particle acceleration is essentially non-thermal, the ICMs may apply to both geometries.

If the electron cooling is dominated by synchrotron self-Compton ($l_\gamma/l_{nth} \ll 1/5$ from the non-magnetic models), the fits suggest that the magnetic field strength is roughly in equipartition with the radiation field. From a qualitative fit of the canonical hard state spectrum of Cygnus X-1, MB09 constrained the magnetic energy density to be strictly below equipartition ($l_B/l_{B_0} < 0.3$). As a consequence, even if our results are consistent with the results obtained for Cygnus X-1, the magnetic field could play a more important role in the physics of GX 339–4, at least in the cutoff state.

In both magnetic models, a transition between the two spectra can be explained by the variations of only two parameters. In any case, the magnetic compactness $l_B$ needs to be variable, along with either the injection slope (in the ICM1) or the maximum energy of the accelerated particles (in the ICM2). All other fit parameters are found to remain constant within the 90% confidence errors. To reproduce the cutoff to excess transition with the ICM1, the magnetic field has to decrease by a factor of 5.4 and the injection slope drops from $\Gamma_{inj} = 3.5$ to 2.5, while in the ICM2, the magnetic field decreases by a factor of 3.2 and the maximum energy jumps from $\gamma_{max} = 15.2$ to 187.

Regardless of the precise acceleration process involved in the accretion flow, the inferred change in the magnetic field strength is expected to have an impact on the cutoff energy of the accelerated electron distribution. Indeed, the maximum energy of the particles is achieved when the energy losses become larger than the gains. In the magnetized models investigated in this paper, the losses at high energies are dominated by synchrotron cooling (Compton and Coulomb cooling are significantly smaller), which obviously depends on the magnetic field strength. In the framework of diffusive shock acceleration for instance, it has been shown that when synchrotron losses are included, the maximum Lorentz factor of the accelerated relativistic electrons satisfies $\gamma_{max} \propto 1/(KB^2)$, where $K$ is the diffusion coefficient (see e.g., Webb et al. 1984 or Marcowith & Kirk 1999). If $K$ is constant, this predicts a maximum energy of $\gamma_{max} \propto 1/B^2$. However, depending on the assumptions made to describe the acceleration mechanism, the diffusion coefficient can depend both on the particle energy and the magnetic field strength: $K(\gamma, B)$. In the frequently used Bohm limit, it is assumed that $K \propto \gamma/B$, implying that the maximum Lorentz factor of the accelerated particles is expected to follow $\gamma_{max} \propto B^{-1/2}$. Although our results are not consistent with the Bohm predictions, the overall behavior remains that the cutoff energy decreases with an increasing magnetic field strength. Moreover, other acceleration mechanisms, such as magnetic reconnection for instance, may give different predictions.

As mentioned in the previous section, the injection index is not universal and may undergo variations in response to chang-
Figure 8. Photon spectra (left panel) and particle distributions (right panel) in the cutoff and excess states, obtained with the BELM code for the best-fit parameters of the ICM2. The photon spectra include neither absorption nor reflection. $\tau(p)$ is the Thomson optical depth per unit momentum. For each particle distribution, the vertical dotted lines show the energy chosen to define the thermal and the non-thermal populations (cf. Section 5.3).

Figure 9. Photon spectra obtained with the BELM code for the best-fit parameters of the ICM2 in the cutoff state (left panel) and the excess state (right panel) with contributions from the various radiation processes: self-absorbed synchrotron emission (green lines), Comptonization off the thermal and non-thermal populations (red and orange lines, respectively), and pair production/annihilation (blue lines). Solid and dashed lines show positive and negative contributions to the spectra, respectively.

The model photon spectra are compared with the contributions from the various radiation processes. Although the lower energy part of the model photon spectra is similar in shape, the underlying electron distributions are quite different (cf. Figure 8). Indeed, the inferred variations of $B$ and $\gamma_{\text{max}}$ not only change the non-thermal part of the distribution but also the temperature of the thermalized component. Figure 9 compares the contributions from the various radiation mechanisms to the model photon spectra, showing separately the Comptonization off thermal and non-thermal particles. As expected, the parameter variations strongly increase the ratio of the non-thermal to the thermal component, resulting in the production of the high-energy tail.

The low energy part of the particle distributions were fitted with a Maxwellian and particles of energy one order of magnitude larger than the inferred temperature were considered to be non-thermal.
In conclusion, the ICM2 provides a framework for a simple, physically motivated interpretation of the data, showing that the spectral variability could be triggered by the variation of only one single parameter, namely the magnetic field strength. Assuming a spherical medium of radius $R = 30 R_G$, a black hole mass of $M = 13 M_\odot$ and a distance of $d = 8$ kpc, we infer a factor of 3 variation of the magnetic field strength, between $B \lesssim 3.9^{+1.5}_{-1.3} \times 10^6 \, G$ and $B \lesssim 1.25^{+0.50}_{-0.45} \times 10^7 \, G$. If the fitted average spectra provide a good estimate of the individual spectra from the single science windows, the timescale of this evolution is at most of the order of hours. Using the standard $\alpha-$prescription (Shakura & Sunyaev 1973), the viscous timescale of the accretion disk at a radius $R$ is given by $t_{\text{visc}} = \alpha^{-1} (H/R)^2 t_K$, where $H/R$ is the aspect ratio of the disk and $t_K$ is the Keplerian period. Using the lower limit $\alpha \gtrsim 0.01$ (quiescent disk) and $H/R \simeq 0.1$ (thin disk), we obtain $t_{\text{visc}} < 10$ minutes for the typical source size $R = 30 R_G$.

The viscous timescale of ADAF-like models is much shorter. Therefore, even if the geometry of the X-ray-emitting region and its dynamical evolution remain uncertain, global changes on timescales of hours are not unrealistic.

6. SUMMARY AND CONCLUSION

We presented an analysis of the high-energy emission of GX 339−4 in a luminous hard state. With respect to the standard cutoff shape of the hard state spectrum, the 25−500 keV INTEGRAL/SPI data revealed the appearance of a variable high-energy excess. The intensity of this hard excess seems to be positively correlated with the total X-ray luminosity and the associated timescale is shorter than 7 hr. We explored the possible physical origins of this variability through an extensive analysis of two averaged spectra, one showing the typical cutoff shape and one showing this prominent high-energy excess.

We used simultaneous RXTE/PCA data to extend the spectral coverage down to 4 keV and fitted the broadband spectra with a variety of physical Comptonization models. Models involving only thermal heating can be ruled out since they are not able to reproduce the high-energy tail. This feature thus confirms that in luminous hard states, the Comptonizing plasma of GX 339−4 contains a fraction of non-thermal particles. Models involving only non-thermal electron acceleration, on the other hand, showed that the thermal part of the spectrum can be produced by an initially non-thermal (power law) distribution which rapidly thermalizes under the effects of synchrotron self-absorption and/or $e^+e^-$ Coulomb collisions.

The relatively good signal-to-noise ratio of the high-energy channels (>150 keV) allowed us to derive meaningful constraints to the model parameters. Depending on the nature of the plasma (magnetized or not), the transition between the two averaged spectra requires the variations of at least two parameters. We found that a magnetized medium subject to non-thermal electron acceleration provides the framework for a simple and physically consistent interpretation of the data. Indeed, the spectral variability could be triggered by only the variations of the magnetic field, implying a subsequent variation of the maximum energy of the accelerated particles.

The quantitative constraints derived from this model yield a very conservative upper limit on the average magnetic field strength in the Comptonizing plasma and show that in the bright hard state, the magnetic energy density could reach equipartition with the radiative energy density.

In conclusion, the presented results suggest that magnetic processes are likely to play a crucial role in the production of the high-energy emission of GX 339−4. We showed that in luminous hard states, the Comptonized emission could originate from a magnetized corona essentially powered through non-thermal particle acceleration, similarly to what is believed to happen in soft states. The hard X-ray emission in both states may therefore be the consequence of a common physical phenomenon.

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No. 2, 2010 VARIABILITY AND SPECTRAL MODELING OF HARD X-RAY EMISSION OF GX 339−4 1035
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