Polarization Change Due to Fast Winds from Accretion Disks

Andrei M. Beloborodov
Stockholm Observatory, S-13336 Saltsjöbaden, Sweden

Received 1997 December 12; accepted 1998 February 4; published 1998 February 24

Abstract

A fraction of the radiation produced by an accretion disk may be Thomson-scattered by a wind flowing away from the disk. Employing a simple plane-parallel model of the wind, we calculate the polarization of the scattered radiation and find that its sign depends on the wind velocity, \( \beta = v/c \). In the case where \( 0.12 < \beta < 0.78 \), the polarization is parallel to the disk normal; i.e., it is orthogonal to the standard Chandrasekhar polarization expected from accretion disks. The velocity of an \( e^\pm \) wind is likely to saturate near the equilibrium value \( \beta_e \sim 0.5 \), for which the accelerating radiation pressure is balanced by the Compton drag. Then the change of polarization by the wind is most pronounced. This may help to reconcile the standard accretion disk model with the optical polarization observations of nonblazar active galactic nuclei.

Subject headings: accretion, accretion disks — polarization — radiative transfer — scattering

1. Introduction

The standard model of a black hole accretion disk predicts the polarization of the emerging radiation. The radiation gets polarized because of electron scattering that strongly dominates the absorption opacity in luminous accretion disks. Polarization of light emerging from an electron plane-parallel atmosphere was calculated by Chandrasekhar (1960) and Sobolev (1963). According to their results, the polarization should be parallel to the disk surface, with a maximum of \( 11.7\% \) when the disk is viewed edge-on. This simple model, however, disagrees with observations: the optical polarization in nonblazar active galaxies viewed edge-on. This simple model, however, disagrees with observations: the optical polarization in nonblazar active galaxies (AGNs) is typically \( \sim 1\% \) or less, and it tends to be parallel to the radio jet that is presumably perpendicular to the disk (Stockman, Moore, & Angel 1984; Antonucci 1992).

The problem of accretion disk polarization has been discussed by a number of authors (e.g., Gnedin & Silant’ev 1978; Loskutov & Sobolev 1981; Sunyaev & Titarchuk 1985; Phillips & Mészáros 1986; Coleman & Shields 1990; Laor, Netzer, & Piran 1990; Matt, Fabian, & Ross 1993; Kartje 1995; Agol & Blaes 1996). In this Letter, we suggest that the original polarization of the disk can be changed as a result of Thomson scattering in a mildly relativistic wind.

Gas outflows are observed in nonblazar AGNs, and evidence for mildly relativistic bulk velocities has been reported (Leighly et al. 1997). The outflows are seen on scales far exceeding the scale of the central source, \( \sim 10^{13} \) cm, and the gasdynamics in the central region is poorly understood. It is quite probable that there is a wind from the innermost region of an accretion disk where the bulk of the observed radiation originates. The wind may form because of a gas outflow from a corona of the disk. It may also form because of an outflow of electron-positron pairs produced by the gamma rays from the disk. In the latter case, the created pairs cool efficiently by inverse Compton scattering down to the Compton temperature \( kT_C \sim 1–10 \) keV, and the cool, light \( e^\pm \) plasma is then easily pushed away by the radiation pressure.

To investigate the effects that a wind could introduce in the pattern of polarization, we consider the simplest plane-parallel model in which the disk is replaced by a homogeneously emitting plane and the wind is assumed to be a cold vertical outflow from the plane (by “cold” we mean that the thermal motions of particles in the outflow are slow compared with the bulk motion). Thomson scattering of frequency-integrated polarized radiation is governed by the transfer equations written down in § 2. The equations account for the relativistic aberration of light in the wind rest frame, which is crucial for the polarization due to scattering. In §§ 3 and 4, we discuss in turn the optically thin and optically thick winds. The results are summarized in § 5.

2. Basic Equations

In a plane-parallel slab, the polarized radiation is represented by the frequency-integrated intensities \( I_l \) and \( I_r \), where the index \( l \) refers to the radiation polarized in the meridional plane (defined by the normal and the ray) and the index \( r \) refers to the polarization perpendicular to this plane. The radiation is axi-symmetric with respect to the normal, with the intensity being a function of \( \mu = \cos \alpha \), where \( \alpha \) is the angle between the ray and the normal. Stationary Thomson transfer in a cold medium moving in the vertical direction with a velocity \( \beta = v/c \) is described by the equations

\[
\frac{\partial I}{\partial \tau} = (1 - \beta \mu) (I - S), \quad \frac{\partial Q}{\partial \tau} = (1 - \beta \mu) (Q - R). \tag{1}
\]

Here \( \tau \) is the scattering optical depth that the medium would have at rest, \( I = I_l + I_r \), and \( Q = I_l - I_r \). \( S \) and \( R \) are source functions representing the scattered radiation.

A simple way to derive the source functions is to transform the radiation field into the wind rest frame where Thomson scattering is coherent. Polarization is invariant with respect to Lorentz boosts along the normal, and the frequency-integrated intensities \( I \) and \( Q \) transform as (see, e.g., Rybicki & Lightman 1979)

\[
I^* (\mu_c) = D^4 I (\mu), \quad Q^* (\mu_c) = D^4 Q (\mu),
\]

where the index \( c \) stands for the comoving frame, \( D = \frac{\mu - \beta}{1 - \beta \mu} \).

\(^1\) Also at Astro-Space Center of P. N. Lebedev Physical Institute, Profsoyuznaya ul. 84/32, 117810 Moscow, Russia.
The transform in exactly the same way. Note that \( I_0 - Q_0 \approx Q_0 - Q_2 \).

In a two-dimensional approach, an initially vertical wind becomes quasi-radial at some height, and its opacity falls off. We consider the simplified one-dimensional problem in which the wind is replaced by a slab of vertically outflowing plasma of a given optical depth \( \tau_0 \). At the outer boundary, \( \tau = 0 \), we assume free escape of the radiation. At the inner boundary, \( \tau = \tau_0 \), we assume the existence of a source of limb-darkened, positively polarized radiation (Chandrasekhar 1960) for which \( I_1 \approx 0.581I_0, I_2 \approx 0.414I_0 \).

3. OPTICALLY THIN WIND

In the case of a small optical depth, \( \tau_0 \), the optically thin approximation holds at angles satisfying the condition, \( \tau_0 = \tau_0(1 - \beta \mu)/\mu < 1 \), and equations (1) yield the scattered radiation emerging from the slab as

\[
I_w \approx \tau_0 S(\mu), \quad Q_w \approx \tau_0 R(\mu).
\]

The polarization degree of the scattered radiation equals

\[
p_w(\mu) = Q_w/I_w = R/S.
\]

From equations (2) and (3), we get

\[
p_w(\mu) = \frac{3(1 - \mu^2)}{8D^3} \frac{\eta}{(3\mu_e^2 - 1)\eta},
\]

where \( \eta = \frac{I_1}{I_0} \approx \frac{I_2}{I_0} \approx \frac{1}{3} \).

The contribution to \( \eta \) from \( Q_0 \) and \( Q_2 \) is \( \sim 30 \) times smaller than that from \( I_1 - I_0/3 \) and can be neglected. The parameter \( \frac{1}{2} < \eta < \frac{1}{3} \) shows to what extent the radiation field is stretched along the normal in the rest frame of the wind, which is of crucial importance for the polarization of the scattered radiation since \( p_w(\mu) \) has the same sign as \( \eta \) at all \( \mu \). In Figure 1, we show the dependence of \( \eta \) on the wind velocity \( \beta \); \( \eta \) changes sign at \( \beta = 0.12 \) and \( \beta \approx 0.78 \). It follows that the scattered radiation remains positively polarized if the wind is slow, \( \beta < \beta_c \). The change in the sign of the polarization happens if the wind has a velocity in the range \( \beta_c < \beta < \beta_2 \).

In the case where the outflow is composed of a light \( e^\pm \) plasma, the timescale to achieve Compton equilibrium with the radiation field is small, and the wind velocity adjusts so that the radiation pressure is balanced by the Compton drag (see, e.g., Phinney 1982). It means that, in the comoving frame, the net radiation flux vanishes, i.e., \( 4\pi I_1 = 0 \), and the equation for the equilibrium velocity, \( \beta_c \), reads

\[
\beta_c = (I_0 + I_3)/(1 + \beta_2^2),
\]

which yields \( \beta_c \approx 0.52 \) for the limb-darkened radiation from the disk. One can check that \( \eta(\beta) \) has a minimum at \( \beta = \beta_c \).

It follows that a wind produces the strongest negative polarization if it is in Compton equilibrium with the radiation field. This fact is illustrated by the diagram embedded in Figure 1 that shows the angular distribution of the disk radiation in the rest frame of the wind. A special feature of the aberration of mildly relativistic \( \beta \) is that the radiation field is effectively “compressed” in the vertical direction. It is this compression that results in the parallel polarization of the scattered radiation, and the compression is naturally strongest at \( \beta = \beta_c \), for which the net flux in the comoving frame vanishes.

\[
\gamma(1 - \beta \mu) \text{ is the Doppler factor, and } \gamma \text{ is the Lorentz factor. When viewed from the comoving frame, the scattered radiation is represented by the source functions, } S(\mu) \text{ and } R(\mu), \text{ given by Sobolev (1963). Transforming the scattered radiation back into the laboratory frame, } S(\mu) = D^{-4}S'(\mu_c) \text{ and } R(\mu) = D^{-4}R'(\mu_c), \text{ we get}
\]

\[
S(\mu) = \frac{1}{D^4} \left[ I_0^b + \frac{1}{\gamma} (3\mu^2_c - 1) \left( I_0^c - \frac{1}{2}I_0^c + Q_0^c - Q_3^c \right) \right],
\]

\[
R(\mu) = \frac{9}{8D^3} \left( I_0^c - \frac{1}{2}I_0^c + Q_0^c - Q_3^c \right),
\]

where \( I_m^c \) and \( Q_m^c \) (\( m = 0, 1, 2 \)) are the moments of the radiation field in the comoving frame,

\[
I_m^c = \int_{-1}^{1} F(\mu_c) \mu_c^m d\mu_c, \quad Q_m^c = \int_{-1}^{1} Q(\mu_c) \mu_c^m d\mu_c.
\]

The \( I_m^c \) are related to the corresponding moments in the lab frame, \( I_m^b \), by

\[
I_0^b = \gamma^2 (I_0 - 2\beta I_1 + \beta^2 I_2), \quad I_1^b = \gamma^2 (\beta^2 I_0 - 2\beta I_1 + I_2),
\]

\[
I_2^b = \gamma^2 (\beta I_0 + I_2) + (1 + \beta^2) I_0.
\]

The \( Q_m \) transform in exactly the same way. Note that \( Q_0^c - Q_0^c = Q_0 - Q_2 \).
Whatever the velocity of the wind, $\beta$, the scattered radiation has an extremum of polarization at $\mu = \beta$. It corresponds to $\mu_\perp = 0$, i.e., to the case where the ray is perpendicular to the normal in the comoving frame, and we denote the extremum as $p_\perp^e$. In Figure 1, $p_\perp^e$ is plotted against $\beta$. For a mildly relativistic wind, $p_\perp^e \approx \eta$, e.g., $p_\perp^e = -19\%$ for $\beta = \beta_\perp$. In the ultra-relativistic limit, the aberrated radiation gets concentrated in a “head-on” direction and $p_\perp^e \rightarrow 100\%$, as discussed by Begelman & Sikora (1987).

An observer viewing a disk with an optically thin outflow will see the sum of unscattered and once-scattered radiation. In Figure 2, we plot the observed polarization as a function of the disk inclination, $\mu$, for the case $\tau_0 = 0.1$ and $\beta = \beta_\perp = 0.52$. The optical depth is small, and the scattering in the wind weakly affects the intensity of the observed radiation. One can see, however, that the wind changes the pattern of polarization. This happens because the polarization extremum of the scattered radiation is boosted into angles $\mu \sim 0.5$, where the original disk polarization is small. Note that the calculated slab model fails at small $\mu$ as $\tau_\parallel$ increases proportionally to $\mu^{-1}$, but the single-scattering approximation is not applicable. In a real wind, the scattering region has a typical height comparable to its radius, and a detailed two-dimensional model is needed to get an exact pattern of polarization. To illustrate the importance of the geometry, let us take a cylinder of height $h$ and radius $r = \sqrt{3}h$ instead of a slab. Then, at $\mu < 0.5$, $\tau_\parallel$ changes to $\tau_\parallel = \sqrt{3}\tau_\parallel(1 - \beta\mu)(1 - \mu^2)^{-1/2}$. Combined with the slab source functions, $S$ and $R$, this yields a rough estimate for the resulting polarization shown by the dotted line in Figure 2.

4. OPTICALLY THICK WIND

An optically thick $e^\pm$ outflow may emerge in the case where a fraction of the disk luminosity is emitted above 511 keV, as hard photons are absorbed by softer X-rays to produce $e^\pm$ pairs (the nonlinear transfer of gamma rays above a luminous disk is discussed in Beloborodov 1998). Compton equilibrium is established in each layer of the outflow, and the bulk velocity is determined by the radiation field itself through the equilibrium condition in equation (6). This means that the velocity should be calculated self-consistently in the transfer problem. Comparing the Compton equilibrium transfer with the classical transfer problem as an electron medium at rest, we note that in addition to a constant net flux, $4\pi I_\parallel(\tau) = \text{const}$, we have $I_\perp(\tau) = \text{const}$ (one can check this by combining the first moment of eq. [1] with eq. [6]). $4\pi I_\parallel/c$ equals the radiation pressure in the vertical direction, and its constancy reflects the condition that the net radiative force acting on $e^\pm$ vanishes in the limit $\beta \rightarrow \beta_\perp$. The equilibrium velocity varies with height, and the pairs keep $\beta \approx \beta_\perp$, since $\beta_\perp$ is a strong attractor in phase space. We have simulated this transfer problem numerically. We take as initial conditions that the disk radiation is propagating freely in the slab, then “switch on” the scattering by the flowing pairs at $\tau = 0$, and follow the evolution of the system until a stationary solution is established. In the calculations, we use a grid that is homogeneous in the $\cos^{-1} \mu$ and $\tau$ directions. The number of grid points is $N_\tau \times N_\mu = 200 \times 300$. The stationary solution, $I(\tau, \mu)$ and $Q(\tau, \mu)$, that we are looking for depends only on the parameter $\tau_0$ and does not depend on the $e^\pm$ density profile, $n(\tau)$. Therefore, we choose the simplest profile, $n(\tau) = \text{const}$. The chosen time step equals $\Delta \tau = 0.5\tau_0/\ln(\eta c N_\tau)$.

The stationary solution for $\tau_0 = 3$ is shown in Figures 3 and 4. Figure 3 displays the wind velocity. In the lowest layers of the slab, the equilibrium velocity is less than $0.5c$ since, besides the disk radiation, there is some flux in the opposite direction as a result of the radiation being backscattered by the outflow. In higher layers, the radiation gets collimated by the scattering on the moving pairs, and the equilibrium velocity increases to $\sim 0.7c$. Figure 4 shows the radiation emerging from the slab: it is significantly anisotropic, and it is negatively polarized with an extremum of $p \approx -14\%$ at $\mu \approx 0.43$.

5. CONCLUSIONS

We conclude that due to Thomson scattering in a wind, the disk radiation originally polarized perpendicular to the normal can acquire parallel polarization. This effect arises because of the relativistic aberration of light in the wind rest frame and takes place when the wind velocity is in the range $0.12 < \beta < 0.78$. Being characteristic of mildly relativistic winds, the production of parallel polarization does not occur in the previously discussed scattering by highly relativistic jets (Begel-
Fig. 4.—Radiation emerging from an optically thick $e^\pm$ wind with $\tau_b = 3$ in comparison with the Chandrasekhar radiation of the disk: (a) the angular dependence of the total intensity (normalized so that the net flux equals $\pi$); (b) the angular dependence of the degree of polarization.

... (Beloborodov & Sikora 1987) that produce strongly beamed radiation with polarization perpendicular to the jet.

An observer viewing a disk through an optically thin wind will see the sum of the radiation from the disk and the radiation scattered in the wind. If the original disk radiation has the perpendicular Chandrasekhar-Sobolev polarization, then the scattering in the wind tends to diminish it. The strongest effect is produced by a wind with velocity $\beta = \beta_c \sim 0.5$, where $\beta_c$ corresponds to a Compton equilibrium of the outflow with the disk radiation field. Such an equilibrium is likely to be established in a wind composed of an $e^\pm$ plasma. We find in this case that a modest optical depth of the wind, $\sim 0.1$, is enough to change the pattern of the disk polarization: the polarization degree is reduced, and there appears a range of inclination angles with negative (parallel) polarization. This change may naturally explain the optical polarization observed in nonblazar AGNs.

We have also found that parallel polarization is produced in an optically thick $e^\pm$ outflow that might form above a black hole accretion disk. The outflow is semirelativistic, and the emerging radiation is markedly beamed along the disk axis. The bulk of the disk radiation gets scattered, its original polarization is lost, and the emerging radiation acquires parallel polarization with an extremum of about $-14\%$ at an inclination angle of $\sim 3\pi/2$. Note that an optically thick $e^\pm$ wind also produces an annihilation feature above $511$ keV, which has a blue-shift up to $\sim 2$, albeit smeared out by the velocity gradient in the wind.

Our general conclusion is that a wind produces parallel polarization when it is near Compton equilibrium with the radiation field. This conclusion is derived for vertical winds in slab geometry, and it remains to be investigated how this conclusion is affected by deviations from this simple model. In particular, a magnetic field may affect the wind direction, especially if magnetic pressure dominates the radiation pressure, and additional relativistic effects are important near Kerr black holes. A detailed two-dimensional model should include a radial dependence of the disk emission and the wind density.

I thank C.-I. Björnsson, J. Poutanen, and R. Svensson for many useful remarks and the referee for helpful comments. I acknowledge support from the Swedish Natural Science Research Council and RFFI grant 97-02-16975.

REFERENCES

Agol, E., & Blaes, O. 1996, MNRAS, 282, 965
Antonucci, R. R. J. 1992, in AIP Conf. Proc. 254, Testing the AGN Paradigm, ed. S. Holt, S. Neff, & C. M. Urry (New York: AIP), 486
Begelman, M. C., & Sikora, M. 1987, ApJ, 322, 650
Beloborodov, A. M. 1998, MNRAS, submitted
Chandrasekhar, S. 1960, Radiative Transfer (New York: Dover)
Coleman, H. H., & Shields, G. A. 1990, ApJ, 363, 415
Gnedin, Y. N., & Silant’ev, N. A. 1978, Soviet Astron., 22, 325
Kartje, J. F. 1995, ApJ, 452, 565
Laor, A., Netzer, H., & Piran, T. 1990, MNRAS, 242, 560
Leighly, K. M., Mushotzky, R. F., Nandra, K., & Forster, K. 1997, ApJ, 489, L25
Loskutov, V. M., & Sobolev, V. V. 1981, Astrofizika, 17, 535
Matt, G., Fabian, A. C., & Ross, R. R. 1993, MNRAS, 264, 839
Phillips, K. C., & Mészáros, P. 1986, ApJ, 310, 284
Phinney, S. 1982, MNRAS, 198, 1109
Rybicki, G. B., & Lightman, A. P. 1979, Radiative Processes in Astrophysics (New York: Wiley)
Sobolev, V. V. 1963, A Treatise on Radiative Transfer (Princeton: Van Nostrand)
Stockman, H. S., Moore, R. L., & Angel, J. R. P. 1984, ApJ, 279, 485
Sunyaev, R. A., & Titarchuk, L. G. 1985, A&A, 143, 374