The Dynamics of Atmospherically Driven Intraseasonal Polar Motion

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ABSTRACT

The atmospheric dynamical processes that drive intraseasonal polar motion are examined with National Centers for Environmental Prediction–National Center for Atmospheric Research reanalysis data and with pole position data from the International Earth Rotation Service. The primary methodology involves the regression of different atmospheric variables against the polar motion excitation function.

A power spectral analysis of the polar motion excitation function finds a statistically significant peak at 10 days. Correlation calculations show that this peak is associated with the 10-day, first antisymmetric, zonal wavenumber 1, normal mode of the atmosphere. A coherency calculation indicates that the atmospheric driving of polar motion is mostly confined to two frequency bands, with periods of 7.5–13 and 13–90 days.

Regressions of surface pressure reveal that the 7.5–13-day band corresponds to the 10-day atmospheric normal mode and the 13–90-day band to a quasi-stationary wave.

The regressions of pole position and the various torques indicate not only that the equatorial bulge torque dominates the mountain and friction torques but also that the driving by the equatorial bulge torque accounts for a substantial fraction of the intraseasonal polar motion. Furthermore, although the 10-day and quasi-stationary wave contributions to the equatorial bulge torque are similar, the response in the pole position is primarily due to the quasi-stationary wave.

Additional calculations of regressed power spectra and meridional heat fluxes indicate that the atmospheric wave pattern that drives polar motion is itself excited by synoptic-scale eddies. Regressions of pole position with separate torques from either hemisphere show that most of the pole displacement arises from the equatorial bulge torque from the winter hemisphere. Together with the above findings on wave–wave interactions, these results suggest that synoptic-scale eddies in the winter hemisphere excite the quasi-stationary wave, which in turn drives the polar motion through the equatorial bulge torque.

1. Introduction

Interest in the atmospheric driving of polar motion, the movement in the position of the pole relative to the earth’s crust, has increased during the past two decades. Part of the motivation for this interest lies in the observation that there is a strong link between polar motion and fluctuations in the amplitude of the two equatorial components of the atmospheric angular momentum (hereafter EAAM) vector. Studies such as those by Eubanks et al. (1988), Chao (1993), and Nastula (1995) have shown that on both interannual and intraseasonal time scales, the linear correlation between EAAM and polar motion, which is a measure of the equatorial angular momentum of the solid earth, has a value that is typically near 0.6. Because of angular momentum conservation within the solid earth–ocean–atmosphere system, excluding tidal influences, the value of this correlation (which is much less than one) implies that in addition to the atmosphere, the ocean must also play an important role in driving polar motion (e.g., Ponte et al. 1998; Ponte and Stammer 1999; Celaya et al. 1999; Nastula and Ponte 1999). Additional motivation for the study of the atmospheric driving of polar motion lies in its importance for the predictability of polar motion. Predictions of pole position are used for satellite navigation and positioning because the tracking of ar-
tificial satellites requires a very accurate knowledge of the location of the satellites relative to the earth’s surface.

Intraseasonal polar motion has been difficult to observe in unfiltered pole position data because it is usually overwhelmed by the much larger-amplitude Chandler and annual wobbles. The Chandler wobble is a free mode of oscillation of the solid earth, which arises because of the oblateness of the solid earth, that is, the polar radius of the earth being shorter than its equatorial radius. The Chandler wobble takes place when the earth’s axis of rotation is not aligned along one of its principal axes. The period of the Chandler wobble is approximately 435 days. The annual wobble, in contrast, is a forced wobble which arises from seasonal shifts in the atmospheric and oceanic mass distribution. The amplitude of these two wobbles is very similar. When the Chandler and annual wobbles constructively interfere, together they can displace the pole position by as much as 10 m. In contrast, when they destructively interfere, which occurs once every 6.4 yr (a period that corresponds to the beat frequency), the intraseasonal displacements of pole position become visible in unfiltered pole position data. The most recent episode of destructive interference took place during the time period from November 2005 through February 2006. Throughout this relatively short time period, intraseasonal polar motion was very clearly observed (Lambert et al. 2006). Figure 1 shows a polar plot of the pole position for the time interval examined by Lambert et al. (2006). In Fig. 1, the axes are aligned along the Greenwich (0°) and 90°E meridians, where the units along both axes are in milliarcseconds (mas), where 1.0 mas ~ 3.0 cm. As is discussed by Lambert et al. (2006), and as can be seen in Fig. 1, there are small-amplitude “loops” in the pole position. These loops take place over a period of about 1–2 weeks, with pole displacements of about 10 cm. In this study, as we will find, atmospherically driven, intraseasonal loops of polar motion are a regularly observed feature of polar motion when the pole position data are filtered to retain only intraseasonal time scales. Our focus in this investigation will be on examining the atmospheric dynamical processes that drive these intraseasonal displacements in the position of the pole.

When studying polar motion, it is convenient to utilize the complex, two-dimensional, polar motion excitation function \( X_p \) = \( X_{p,1} + iX_{p,2} \), which is defined as (see Lambeck 1980; Barnes et al. 1983)

\[
X_p = m + (i/\alpha_0)dm/dt, \tag{1}
\]

where \( m = m_1 + im_2 \) is the complex angular distance, as illustrated in Fig. 1. The subscripts 1 and 2 in \( X_p \) and \( m \) denote the components aligned with the 0° and 90°E meridians, respectively; \( \alpha_0 \) is the complex Chandler frequency, given by \( \alpha_0 = 2\pi(1 - i/2Q)/T_0 \), where \( T_0 \) is the Chandler period; \( Q \) is a dissipation factor; \( t \) is time; and \( i = \sqrt{-1} \). For this study, we set \( T_0 = 435 \) solar days and \( Q = 100 \).

Most studies that link polar motion to atmospheric driving utilize the two-dimensional atmospheric excitation function \( \chi_a \), a quantity which is closely related to the EAAM vector. If fluctuations in the equatorial oceanic angular momentum vector are not considered, then

\[
m + (i/\alpha_0)dm/dt = \chi_p = \chi_a - (i/\Omega)d\chi_a/dt, \tag{2}
\]

where \( \Omega \) the angular velocity of the solid earth (see Barnes et al. 1983). The relationship between \( X_p \) and \( \chi_a \) can be simplified by noting that the first term on the right-hand side (rhs) of (2) is about one order of magnitude larger than the second term (Barnes et al. 1983), implying that \( X_p \approx \chi_a \). In a coordinate system fixed with the rotating earth, the 0° and 90°E components of \( \chi_a \) are as follows:

![Fig. 1. A polar plot of the pole position for the time period from 1 Nov 2005 through 14 Feb 2006. This is the same time period examined by Lambert et al. (2006), when the Chandler and annual wobbles destructively interfere. The abscissa (x axis) is aligned along the 0° meridian and the ordinate (y axis) is aligned along the 90°E meridian. The units on the axes are milliarcseconds (mas). This plot was calculated on the IERS Web site (http://iers.obspm.fr).](image-url)
where $p_s$ is the surface pressure; $u$ and $v$ are the zonal and meridional winds relative to the earth’s surface, respectively; and $\lambda$, $\phi$, and $p$ are the longitude, latitude, and pressure, respectively. [See Barnes et al. (1983), Eubanks (1993), and Zhou et al. (2006), which present updated values for the constants in $\chi_{a1}$.] The quantities $A$ and $C$ are the polar and equatorial principal moments of inertia of the solid earth, respectively; $R$ is the earth’s radius; and $g$ is the gravitational acceleration. The coefficients 1.098 and 1.591 3 in (3) and (4) correspond to adjustments that arise from the sum of the influences of rotational and surface loading (Barnes et al. 1983). If the earth were perfectly rigid, then both coefficients would be equal to unity, and the atmospheric excitation function would be exactly proportional to the EAAM vector. (The mathematical expression for the EAAM vector is included in the appendix.)

An approximate equation relating $\chi_p$ to the atmospheric torques can be obtained by noting that the EAAM tendency in an inertial coordinate system, which is equal to the sum of the torques applied to the atmosphere, is very similar to the tendency of the atmospheric excitation function also in an inertial coordinate system; that is,

$$\chi_a - (i/\Omega) d\chi_a/dt \approx 1/(\Omega(C-A)) [EM - (i/\Omega)d(EM)/dt]$$

$$= -i/(\Omega^2(C-A))(BT + MT + FT),$$

where $EM$, $BT$, $MT$, and $FT$ are the two-dimensional, complex representations of the EAAM vector, equatorial bulge torque, mountain torque, and friction torque, respectively. (The mathematical expression for these three torques can be found in the appendix.) Using (2), this allows us to express $\chi_p$ in terms of the torques; that is,

$$\chi_p \approx -i/(\Omega^2(C-A))(BT + MT + FT). \quad (5)$$

In this study, to examine the atmospheric processes that drive intraseasonal polar motion, we will regress these torques and other atmospheric variables against the two components of $\chi_p$. On intraseasonal time scales, the second term on the rhs of (1) is about one order of magnitude larger than the first term (on interannual time scales, these two terms are of similar magnitude). Therefore, for intraseasonal time scales, regression against $\chi_p$ corresponds very closely with regression against the time derivative of the pole position $dm/dt$. With regard to the torques, this indicates that the pole position $m$ depends on the time integral of the torques and not on the torques themselves. An implication of this relationship between the pole position and torques can be seen by first writing the pole position $m$ and the sum of the three torques $\tau$ in spectral form; that is,

$$m = \sum_\omega m_\omega e^{i\omega t}; \quad \tau = \sum_\omega \tau_\omega e^{i\omega t}.$$  

Then, after using the approximation $\chi_p = dm/dt$, and substituting the above spectral representations for $m$ and $\tau$ into (5), we can rewrite (5) as

$$\omega m_\omega = -\tau_\omega/[\Omega^2(C-A)]. \quad (6)$$

The implication of (6) is that the size of the pole position displacements depends on both the magnitude of the torques and their frequency distribution, with lower frequencies driving the largest changes in pole position. There are many outstanding questions on the subject of atmospherically driven, intraseasonal polar motion. For example, several studies have shown that intraseasonal EAAM is related to the 10-day westward propagation (Brzezinski 1987; Egger and Hoinka 2002; Feldstein 2003b, 2006) of the first antisymmetric, zonal wavenumber 1, Rossby wave normal mode. In Feldstein (2006), it was shown that fluctuations in the amplitude of the EAAM vector are influenced by equatorial mixed Rossby–gravity waves, wave–zonal mean flow interaction, and wave–wave interactions in the midlatitude upper troposphere. However, it is an open
question whether these processes influence intraseasonal polar motion. There is also the question of which atmospheric torques dominate the driving of polar motion. Even though the equatorial bulge torque has been shown to be the largest torque in the intraseasonal EAAM budget (Egger and Hoinka 2002; Feldstein 2006) because changes in pole position depend upon the frequency distribution of the torques, it is not necessarily the case that the equatorial bulge torque also dominates intraseasonal polar motion. Another question is which atmospheric time scales drive polar motion. As we will find, two particular frequency bands are the main atmospheric drivers of intraseasonal polar motion. One of these bands will be shown to correspond to the 10-day atmospheric normal mode, as mentioned above, and the other band to a quasi-stationary wave. This result leads to the question of what dynamical processes excite these two waves. Another interesting question is how these two waves are related to the climatological stationary eddies and midlatitude baroclinicity. For example, do the amplitudes of these two waves, and thus the torques and the ensuing polar motion, depend on the amplitude of the climatological stationary eddies and the climatological midlatitude baroclinicity? In this study, each of the above questions about the atmospheric driving of intraseasonal polar motion will be investigated.

The methodology is described in section 2, followed by an examination of the power spectra, normal modes, and coherency in section 3. The EAAM budget is presented in section 4, and a calculation of the atmospheric wave fields associated with polar motion is shown in section 5. Section 6 examines the response of the pole position to the atmospheric torques, and in section 7 the dynamical mechanisms that excite the atmospheric wave fields are investigated. The conclusions are given in section 8.

2. Data and methodology

The daily pole position data used in this study, which were obtained from the International Earth Rotation Service (IERS), span the years from 1962 through 2005. To examine the atmospheric contribution to the driving of polar motion, daily (0000 UTC) National Centers for Environmental Prediction–National Center for Atmospheric Research (NCEP–NCAR) reanalysis data are used. The atmospheric variables used in this study include surface pressure, 300-hPa streamfunction, and the equatorial bulge, friction, and mountain torques. Polar motion is examined both during the boreal winter (December–February) and the austral winter (June–August). To examine the role of tropical convection, we use outgoing longwave radiation (OLR) data, which are produced by the National Oceanic and Atmospheric Administration (NOAA). The OLR data used span the years 1979 through 1995. Except for the calculations of mountain torque and OLR, all data are truncated at rhomboidal 30 resolution. For the mountain torque, the calculation is performed on the NCEP–NCAR triangular 62 grid, and for the OLR, the analysis is performed on a 2.5° latitude × 2.5° longitude grid. A 90-day high-pass digital filter is applied to all variables. This filter removes both the seasonal cycle in the atmospheric data and the Chandler and annual wobbles from the pole position data. As discussed in the introduction, the primary methodology used in this study is linear regression of different atmospheric variables against the two components of $\chi_p$. All calculations will correspond to a one standard deviation anomaly in both $\chi_p$ components.

In this study, for the calculation of the equatorial bulge torque (see the appendix), the observed surface pressure values at all grid points over the earth’s land and ocean surface are used. This common representation of the equatorial bulge torque implies that the oceanic mass does not respond to fluctuations in atmospheric pressure. This contrasts with another approach used in many studies in which the ocean is treated as an inverted barometer (IB), whereby the pressure at the bottom of the ocean completely adjusts to changes in atmospheric pressure. In a recent study, with highly accurate pole position data, Lambert et al. (2006) found that during some time periods the non-IB approach (as adopted in this study) was more accurate and at other times the IB correction yielded more precise solutions. As discussed in Lambert et al. (2006), these findings imply that the actual oceanic response to atmospheric pressure fluctuations lies between those of the non-IB and IB approaches.

The unfiltered and 90-day high-pass pole position time series are shown for the 0° and 90°E meridians in Fig. 2. As can be seen, the unfiltered time series are dominated by the annual and Chandler wobbles and a negative secular trend (Fig. 2b), with the intraseasonal polar motion comprising just a small fraction of the unfiltered variance. (The secular trend in Fig. 2b, which has been observed for more than 100 yr, has not yet been fully explained.) Because the periods of the annual and Chandler wobbles are similar, the sum of these two wobbles takes the form of beats, as can be seen in Fig. 2. The correlation between the corresponding components of the 90-day high-pass $\chi_p$ and $\chi_a$ time series is calculated separately for each of the years between 1962 and 2005. As can be seen in Fig. 3, beginning about 1982 there was a marked increase in the
correlations between these two quantities, from values close to zero up to values near 0.6 and 0.7. [Very similar findings were noted by Nastula and Salstein (1999), including a marked increase in the correlations beginning about 1983.] Assuming that the EAAM data have been of high quality well before 1982, these results indicate a substantial improvement in pole position data on intraseasonal time scales after 1982. Therefore, for this study, we will limit our examination of polar motion to the years from 1982 to 2005.

Because of the dominance of the Chandler and annual wobbles in Fig. 2, the pole position values at the beginning and the end of the time series differ substantially. As a result, the 90-day high-pass filter distorts the first and last 100 days of the dataset. To deal with this problem, after the 90-day high-pass filter is applied, the pole position data for the years 1982 and 2005 are discarded. As a result, the data examined in this study only cover the years 1983 through 2004.

3. Power spectra, normal modes, and coherency

The power spectra for both components of the polar motion excitation function $\chi_p$ are illustrated in Fig. 4. These spectra were obtained by calculating the individual power spectra separately for each boreal and austral winter season and then averaging. To examine statistical significance, each frame in Fig. 4 also illustrates the power spectrum for a first-order autoregressive process, along with the corresponding 95% a priori and a posteriori confidence levels. The number of degrees of freedom for these confidence level curves was specified as twice the number of seasons, that is, 2 degrees of freedom for each seasonal periodogram value. This specification for the number of degrees of freedom is based on the assumption that successive wintertime series are uncorrelated. As can be seen in Fig. 4, all power spectra have statistically significant peaks centered at or near 10 days (Fig. 4c). A spectral peak at this
period is perhaps to be expected, given the strong 10-day peak in the EAAM power spectra, as discussed in the introduction. The other spectral peaks in Fig. 4, centered near 7 and 15–20 days, are also found in EAAM power spectra (see Fig. 3 in Feldstein 2006).

To examine whether the 10-day spectral peaks in Fig. 4 are indeed related to the 10-day normal mode in the atmosphere, we follow a procedure very similar to that used in Feldstein (2003b, 2006), which involves the normal mode solutions to the shallow water equations on the sphere (Eubanks et al. 1988; Bell 1994; Egger and Hoinka 1999). These solutions correspond to the orthonormal Hough modes of Laplace’s tidal equations (Longuet-Higgins 1968; Kasahara 1976, 1980). We use the same notation as in Feldstein (2003b, 2006), where $k = 1$ denotes zonal wavenumber 1 and $n = 2$ indicates the first antisymmetric mode. The equivalent depth is set to 5.8 km. For both the boreal and austral winters, new time series are generated by projecting the daily surface pressure field onto the $k = 1, n = 2$ Hough mode. Two time series are generated, for the 0° and 90°E meridians. These two time series, which both have strong 10-day spectral peaks (Feldstein 2006), together represent the $k = 1, n = 2$ Hough mode contribution to the surface pressure field. As shown in Feldstein (2006), the 10-day spectral peaks for these two time series correspond to westward propagation of the $k = 1, n = 2$ atmospheric normal mode. These two time series are then linearly correlated with the polar motion excitation function time series (see Table 1). Because these correlations are all fairly large and statistically significant well above the 95% confidence level, we can conclude that the 10-day spectral peak in polar motion is associated with the 10-day, $k = 1, n = 2$, atmospheric normal mode.

We next examine the coherency between corresponding components of the polar motion and atmospheric excitation function time series (see Fig. 5, which shows the squared coherency). Along with the coherency, Fig. 5 illustrates the 95% and 99% confidence levels for a zero null hypothesis. As in the power spectral calculations, the coherency was calculated separately for each boreal and austral winter season and then averaged. Figure 5 shows that the coherency exceeds the 99% confidence level at most periods greater than about 8 days. The largest values are seen to occur over a broad range of periods centered near 10 days, with the highest values being observed for the 90°E meridian, during the boreal winter. These results indicate that for periods near 10 days the least squares regression line between the polar motion and atmospheric excitation functions

<table>
<thead>
<tr>
<th>Correlation</th>
<th>$x_p(\text{bor}, 0^\circ)$</th>
<th>0.43</th>
</tr>
</thead>
<tbody>
<tr>
<td>$x_p(\text{bor}, 90^\circ E)$</td>
<td>0.69</td>
<td></td>
</tr>
<tr>
<td>$x_p(\text{aus}, 0^\circ)$</td>
<td>0.42</td>
<td></td>
</tr>
<tr>
<td>$x_p(\text{aus}, 90^\circ E)$</td>
<td>0.50</td>
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accounts for between 50% and 70% of the explained variance, and for longer periods it accounts for between 40% and 60% of the explained variance. For periods less than 8 days, the results are markedly different, revealing coherencies with values that are often below the 95% confidence level, corresponding to only about 30% of the explained variance.

4. Equatorial atmospheric angular momentum budget

To examine the relative strength of the torques that drive intraseasonal polar motion, we regress each term on both sides of the equation for EAAM conservation; that is,

\[
\frac{d\text{EM}}{dt} = \frac{d\text{EM}}{dt} + \Omega \times \text{EAAM} = \text{FT} + \text{MT} + \text{BT},
\]

against the polar motion excitation function. In (7), the subscripts \(i\) and \(r\) denote time derivatives in the inertial and rotating coordinate systems, respectively. The regression coefficients for the equatorial bulge and mountain torques are illustrated in Fig. 6 for the boreal and austral winters for both the 0° and 90°E meridians. As can be seen, for those lags between −5 and +5 days, the EAAM tendency and the equatorial bulge torque closely resemble each other, but the mountain torque is about 5 times smaller than the equatorial bulge torque.
and the friction torque (not shown) is much smaller yet (see Egger and Hoinka 2002; de Viron et al. 1999; Feldstein 2006).

The coherency results of Fig. 5 suggest that the atmospheric driving of intraseasonal polar motion can be separated into three different frequency bands; 2–7.5, 7.5–13, and 13–90 days. The 7.5–13-day band corresponds to the broad peak centered near 10 days in all four frames of Fig. 5. As can be seen, the 7.5-day cutoff period is fairly well defined, whereas the 13-day cutoff is to some extent arbitrary, especially in Fig. 5b. However, as we will find, this choice of a 13-day cutoff period is effective for identifying different dynamical processes. A regression calculation of each of the terms in (7) was also performed separately for these three frequency bands (not shown). Again, it was found that the equatorial bulge torque is dominant, with the mountain torque being much smaller and the friction torque negligible.

5. Atmospheric wave fields

We next examine the atmospheric wave fields associated with polar motion. This is performed by linearly
regressing the anomalous surface pressure field against
the two components of the polar motion excitation
function for both the boreal and austral winters. This
calculation is performed for both the 7.5–13-day and
13–90-day bands.

We focus on the surface pressure field regressed
against the 90°E component of the polar motion excitation
function for the boreal winter because that com-
ponent exhibits the largest equatorial bulge torque
(Fig. 6b). Figure 7 shows the regressed 7.5–13-day sur-
fase pressure field. As can be seen, the lag 0-day surface
pressure field (Fig. 7d) is composed of a zonal wave-
number 1 pattern that is antisymmetric across the equa-
tor, with the extrema being located in the midlatitudes
near the 90°E and 90°W meridians. This is the spatial
pattern that is expected, given that \( \chi_a \) is closely related
to \( \chi_a \) (Fig. 5), and a surface pressure pattern of the form
shown in Fig. 7 maximizes the mass contribution to \( \chi_a \).

As many studies have shown (see Barnes et al. 1983),
the first term on the rhs of (3) and (4) dominates the
atmospheric excitation function.] The surface pressure
field shown in Fig. 7 indicates a wave field that under-
goes westward propagation. By comparing Figs. 7a and
7f, it can be seen that the period of this westward propa-
gation is 10 days, consistent with our findings for the
power spectra (Fig. 4a), coherency (Fig. 5b), and the
normal mode analysis of section 3. A similar regression
calculation was performed for the 13–90-day surface
pressure field (Fig. 8). The results indicate a wave pat-
tern that grows and decays while undergoing very slight
westward propagation. For the rest of this study, wave
fields showing these properties will be referred to as the
10-day and quasi-stationary waves, respectively.

Figures 7 and 8 indicate substantial differences be-
between the winter (Northern Hemisphere) and summer (Southern Hemisphere) hemispheres for the 10-day and quasi-stationary wave fields. Among these differences are the following: (i) the contribution from the winter hemisphere exceeds that from the summer hemisphere, for both the quasi-stationary and 10-day waves; (ii) within the winter hemisphere, the quasi-stationary wave field has a larger amplitude than that of the 10-day wave; and (iii) for the summer hemisphere, it is the 10-day wave that has a larger amplitude than the quasi-stationary wave.

The regressed surface pressure fields were also calculated for the boreal winter, 0° meridian and the 0° and 90°E meridians of the austral winter. Each of these calculations revealed characteristics that were similar to those for the 90°E meridian.

6. Pole position

Having examined the torques and the associated atmospheric circulation patterns, we next investigate the fluctuations in pole position that are driven by the torques. We begin by first showing that the equatorial bulge torque regressed against the polar motion excitation function for the boreal winter, 90°E meridian and the austral winter, 0° meridian. The essential properties of the atmospherically driven polar motion are captured with lagged regressions for these two components. The regressed equatorial bulge torques are illustrated as polar plots (Fig. 9), where the abscissa is the 0° component and the ordinate the 90°E component. The 7.5–13- and 13–90-day bands are shown, along with that for 2–90 days. Each point on the curves in Fig. 9 indicates the magnitude (distance from the origin) and the direction (see the caption for additional detail) of the equatorial bulge torque over the time interval from lag –20 to lag +20 days.

A comparison of Figs. 9b and 9c and Figs. 9e and 9f indicates that the largest equatorial bulge torques for the 7.5–13-day band are comparable to those for 13–90-day band. Furthermore, as expected, Fig. 9 shows...
that the period of the 7.5–13-day equatorial bulge torque is much shorter than that for the 13–90-day band. (The shorter period for the 7.5–13-day torques in Figs. 9b and 9e can be identified by the presence of multiple large-amplitude loops.)

The regression of the pole position \( (m = m_1 + im_2) \) against the same polar motion excitation functions is shown in Fig. 10. The pole position is also displayed in a polar plot. As can be seen, the largest displacement in the pole position is about 3.0 mas or 9.0 cm (Fig. 10a). These values are very close those observed by Lambert et al. (2006) (see Fig. 1) for the time period when the Chandler and annual wobbles destructively interfere. For the boreal winter, 90°E meridian, most of the polar motion is aligned along the same meridian, whereas for the austral winter, 0° meridian, the polar motion has similar contributions from both meridians. In all frames in Fig. 10, although not indicated, the polar motion is retrograde. Figure 10 also indicates that the polar motion associated with the boreal winter, 90°E meridian has about twice the amplitude of that for the austral winter, 0° meridian, which is consistent with the magnitude of the equatorial bulge torques in Fig. 9. The most striking feature of Fig. 10, perhaps, is that the intraseasonal polar motion is completely dominated by the 13–90-day band. The relatively small contribution to the polar motion from the 7.5–13-day band contrasts with Fig. 9, where the driving by the 7.5–13-day and 13–90-day equatorial bulge torques is shown to be similar. A possible explanation for the dominance of the 13–90-day band for polar motion can be attained by referring to (6), where it was shown that the response of the pole position to the driving by atmospheric torques favors lower frequencies.

To address the question of why the changes in pole position are primarily from the 13–90-day band, we use (2) and (5), which together indicate that the change in pole position over a time interval \([t_1, t_2]\), solely from equatorial bulge torque, is equal to the integral of the torque over the same time interval; that is,

\[
\Delta m = \int_{t_1}^{t_2} T(t) \, dt
\]
The pole position is estimated by integrating the regressed equatorial bulge torque, starting from the lag 20-day regressed value of \( m \). The results of this integration are shown in Fig. 11. To be able to address the question of whether the changes in the pole position are primarily a response to the 13–90-day equatorial bulge torque, it is necessary that Figs. 10 and 11 show a reasonable degree of resemblance. Although the match between Figs. 10 and 11 cannot be exact because of the oceanic contribution to the driving of intraseasonal polar motion (Fig. 5), there are still many similarities between Figs. 10 and 11. For example, Fig. 11c captures the primarily rectilinear movement shown in Fig. 10c, Fig. 11f illustrates the more circular movement observed in Fig. 10f, and Figs. 11b and 11e indicate the much shorter time-scale periodicity displayed in Figs. 10b and 10e. Consistent with (6), the extent of these similarities suggests that a substantial fraction of the displacement in pole position is a response to the driving from the 13–90-day, quasi-stationary wave contribution to the equatorial bulge torque.

We also examined the polar motion that is driven by separate winter and summer hemispheric contributions to the equatorial bulge torque. When the summer hemispheric torque was set to zero, the resulting pole position plots (not shown) closely resembled those in Fig. 11. This suggests that the atmospheric driving of intraseasonal polar motion is mostly due to the quasi-stationary waves of the winter hemisphere.

The same methodology was also applied to the mountain torque, where the pole position was obtained by performing a time integration of the regressed mountain torque. It was found that the pole position displacements were very small, much smaller than those for the 7.5–13-day equatorial bulge torque shown in Fig. 11. A possible explanation for this very small polar motion response can again be attained by referring to (6). According to (6), as was discussed, high-frequency torques drive very weak pole displacements. Consistently, a calculation of the power spectra for the mountain torques indicated that a large fraction of the
7. Atmospheric processes

In a recent study (Feldstein 2006), it was shown that the amplitude of the EAAM vector is influenced by the following three processes: (i) poleward propagating Rossby waves that are excited by convection associated with equatorial mixed Rossby–gravity waves, (ii) wave–wave interactions, and (iii) wave–mean flow interaction. We next investigate the extent to which these processes also influence polar motion, using an approach that is similar to that of Feldstein (2006).

The impact of equatorial mixed Rossby–gravity waves was investigated by regressing the OLR and 300-hPa streamfunction fields against the time derivative of $\chi_{p,1}^2 + \chi_{p,2}^2$. In contrast to Feldstein (2006), a clean zonal wavenumber 1 pattern in the tropics was not obtained. This suggests that although equatorial mixed Rossby–gravity waves do influence EAAM, they do not appear to impact polar motion.

To evaluate wave–wave interactions, the zonal wavenumber power spectrum for the 300-hPa geopotential height field was regressed separately against $\chi_{p,1}^2$ and $\chi_{p,2}^2$. The power spectra were evaluated at each day and for each latitude. The results are shown in Fig. 12, where the regressed power spectra are summed over all latitudes. As can be seen, between lag $-5$ and lag 0 days, the power of the shorter-scale eddies, particularly zonal wavenumber 4, declines at the same time as that for zonal wavenumber 1 increases. These results resemble those for EAAM (Feldstein 2006), and are consistent with the hypothesis that the zonal wavenumber 1 field associated with EAAM and polar motion derives part of its energy from shorter-scale waves. In another study (Feldstein 2008, unpublished manuscript), a composite analysis of the streamfunction tendency equation was performed to examine the dynamical mechanisms that excite the zonal wavenumber 1 pattern associated with intraseasonal EAAM variability. Consistent with the findings of this study, that analysis showed that the zonal wavenumber 1 pattern is primarily excited by shorter, synoptic-scale eddies.
The physical mechanism by which the wave–wave interactions alter EAAM involves the excitation of a nonzero, vertically integrated, horizontal divergence at zonal wavenumber 1, which will change the surface pressure field at the same zonal wavenumber and thus also alter the equatorial bulge torque. We illustrate this mathematically with the following derivation for the 0° component of the equatorial bulge torque \[ \text{(A7) in the appendix} \]. Beginning with the continuity equation, \[ \frac{\partial \rho}{\partial t} = -\int_h^\infty \mathbf{\nabla}_H \cdot (\rho \mathbf{v}) \ dz. \] (9)

Next, integrating forward in time from \( t_0 \) to \( t \) yields

\[ p_s(x, y, t) = p_s(x, y, t_0) - \int_{t_0}^t \int_h^\infty \mathbf{\nabla}_H \cdot (\rho \mathbf{v}) \ dz \ d\tau. \] (10)

The substitution of (10) into (A7), combined with the approximation for quasigeostrophic flows that \( \mathbf{\nabla}_H \cdot (\rho \mathbf{v}) = \rho_s(z) \mathbf{\nabla}_H \cdot \mathbf{v} \), where \( \rho_s(z) \) is the time-independent, horizontally averaged density, results in

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**Fig. 12.** The anomalous power spectra regressed against the (a) boreal winter \( X^2_{P,1} \), (b) the boreal winter \( X^2_{P,2} \), (c) the austral winter \( X^2_{P,1} \), and (d) the austral winter \( X^2_{P,2} \). The contour interval is 5 m\(^2\). Solid (dashed) contours are positive (negative), and the zero contour is omitted. Negative values are shaded.
EM_1(t) = EM_1(t_0) + \frac{\Omega R^4}{g} \int_{t_0}^{t} \int_{h}^{\infty} \int_{0}^{2\pi} \int_{-\pi/2}^{\pi/2} \rho_s(z) \nabla_H \cdot \mathbf{v} \sin^2 \phi \cos \lambda d\phi d\lambda dz d\tau.

This equation simply shows that fluctuations in the equatorial bulge torque arise from changes in the zonal meridional heat fluxes showed that the equatorial bulge torque is dominant, with the mountain torque being about 5 times smaller and the friction torque even smaller than that. A regression of the surface pressure field against the polar motion excitation function was performed to determine the atmospheric wave field associated with the equatorial bulge torque. These calculations showed that the 7.5–13-day contribution to the equatorial bulge torque corresponds to the 10-day atmospheric normal mode, and the 13–90-day contribution to a quasi-stationary Rossby wave that has a spatial structure similar to that of the 10-day atmospheric normal mode. These waves were referred to as the 10-day and quasi-stationary waves, respectively.

To determine the impact of the various torques on polar motion, two types of pole position calculations were performed. In one calculation, the pole positions were regressed against the polar motion excitation function. In a second calculation, the regressed torques were integrated forward in time to yield the pole position changes that are driven by the torques. These calculations found that changes in pole position due to atmospheric driving result almost entirely from the equatorial bulge torque. Furthermore, the large degree of similarity between these two pole position calculations indicates that a substantial fraction of the intraseasonal polar motion is due to the equatorial bulge torque. More importantly, even though the contribution to the equatorial bulge torque is similar for the 10-day and quasi-stationary waves, it was found that pole position displacements arise primarily from the quasi-stationary wave driving. This difference in the response of the polar motion to the 10-day and quasi-stationary waves was explained by the fact that low-frequency torques drive larger pole displacements.

Further examination of the regressed power spectra and meridional heat fluxes showed that the equatorial

Equation (13) links the equatorial bulge torque to wave–wave interactions through the nonlinear, relative vorticity advection term (\mathbf{v} \cdot \nabla \xi) in the integrand of (13). As discussed in studies such as Feldstein (2007), horizontal divergence in the midlatitudes can be understood as corresponding to the secondary circulation induced by both relative and planetary vorticity advection. The secondary circulation maintains thermal wind balance (Holton 2004). Therefore, it is this adjustment to thermal wind balance, caused by wave–wave interactions that can excite the equatorial bulge torque.

Wave–zonal mean flow interaction was examined by regressing the zonal wavenumber 1850-hPa, meridional heat flux against \chi_{p,1} and \chi_{p,2}. As in Feldstein (2006), the regressed meridional heat flux is actually equatorward (Fig. 13), indicating that wave–zonal mean flow interaction reduces the amplitude of the zonal wavenumber 1 pattern associated with polar motion. This suggests that synoptic-scale eddies, via the wave–wave interactions, are the primary generator of the atmospheric wave field that in turn drives intraseasonal polar motion.

8. Conclusions

The atmospheric dynamical processes that drive intraseasonal polar motion were investigated. The primary methodology involved the regression of different atmospheric variables against the polar motion excitation function. Both the boreal and austral winters were examined. Power spectral analyses of the polar motion excitation function found several statistically significant peaks, including a 10-day peak, as in several studies of EAAM (Brzezinski 1987; Egger and Hoinka 2002; Feldstein 2003b, 2006). A series of linear correlations indicated that the 10-day spectral peak is associated with the 10-day, first antisymmetric, zonal wavenumber 1, Rossby mode of the atmosphere. A coherency calculation showed that the strongest relationship between the polar motion and atmospheric driving occurs for
bulge torque appears to be excited by wave–wave interactions that involve synoptic-scale waves. Additional calculations of the pole position for the torques from separate hemispheres, for both the boreal and austral winters, indicated that most of the pole displacement was from the winter hemisphere torques. Together with the above result on the role of wave–wave interactions, these findings suggest that the ultimate source of the intraseasonal polar motion are midlatitude synoptic-scale eddies of the winter hemisphere. It appears that these eddies excite the quasi-stationary wave field, which in turn drives the polar motion.

The important role played by quasi-stationary waves of the winter hemisphere alludes to the possibility that some features of the seasonally varying climatological flow may be important for determining the strength of the intraseasonal polar motion. To evaluate the role of the climatological flow, a brief comparison is made between the amplitude of the quasi-stationary and 10-day waves and that of the climatological stationary eddies and midlatitude baroclinicity. For the quasi-stationary waves, we found that these waves are strongest for the NH boreal winter, followed by the SH austral winter, the NH austral winter, and then the SH boreal winter (the amplitude of these waves is measured by the sum of the squares of the 0° and 90°E components). The...
same sequence is observed for the 300-hPa, climatological, stationary eddies. (See the NCEP–NCAR Reanalysis Electronic Atlas from the NOAA/Earth System Research Laboratory Web site, www.cdc.noaa.gov/ncep_reanalysis.) This suggests that the climatological stationary eddies may be playing a role in the generation of the quasi-stationary waves, such as those shown in Fig. 8. The plausibility of this link between the climatological stationary eddies and the quasi-stationary waves is hinted at by some of the properties of large-scale, quasi-stationary, atmospheric teleconnection patterns, such as the North Atlantic Oscillation and the Pacific–North American pattern. These teleconnection patterns are strongest during the NH boreal winter, and their dynamics is closely tied to the presence of large-amplitude climatological stationary eddies in the upper troposphere (Feldstein 2002, 2003a). Such a link may also be related to the findings of Nastula and Salstein (1999), Salstein and Rosen (1989), and Nastula et al. (2007, manuscript submitted to J. Geophys. Res.), who showed that there is a strong longitudinal sensitivity in the covariance between regional EAAM and polar momentum. This is because, as with the teleconnection patterns, the climatological stationary eddies may cause the quasi-stationary waves to attain their largest amplitude at preferred longitudes.

For the 10-day wave, the sequence is the same as that for the quasi-stationary waves of the previous paragraph, except that the amplitude for the SH boreal winter exceeds that for the NH austral winter. This sequence for the 10-day wave matches that of the climatological baroclinicity, as measured by the midlatitude zonal mean potential temperature gradient at 850 hPa (www.cdc.noaa.gov/ncpe_reanalysis). Because of the close relationship between synoptic-scale eddies and baroclinicity (e.g., Blackmon et al. 1984 and references therein), this result suggests that perhaps the 10-day wave is excited by high-frequency, synoptic-scale eddies.

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APPENDIX

Equations for the Torques

In a coordinate system fixed to the rotating earth, the equations for the equatorial bulge, mountain, and friction torques can be written as follows (Bell 1994; Egger and Hoinka 2002):

\[ \text{BT}_1 = \frac{\Omega^2 R^4}{g} \int_{-\pi/2}^{\pi/2} \int_{0}^{2\pi} p_s \sin\phi \cos^2\phi \sin\lambda \, d\phi \, d\lambda, \] (A1)

\[ \text{BT}_2 = -\frac{\Omega^2 R^4}{g} \int_{-\pi/2}^{\pi/2} \int_{0}^{2\pi} p_s \sin\phi \cos^2\phi \cos\lambda \, d\phi \, d\lambda, \] (A2)

\[ \text{MT}_1 = -\int_{-\pi/2}^{\pi/2} \int_{0}^{2\pi} R^2 \left( p_s \frac{\partial h}{\partial \phi} \cos\phi \sin\lambda - p_s \frac{\partial h}{\partial \lambda} \sin\phi \cos\lambda \right) \, d\phi \, d\lambda, \] (A3)

\[ \text{MT}_2 = \int_{-\pi/2}^{\pi/2} \int_{0}^{2\pi} R^2 \left( p_s \frac{\partial h}{\partial \phi} \cos\phi \cos\lambda + p_s \frac{\partial h}{\partial \lambda} \sin\phi \sin\lambda \right) \, d\phi \, d\lambda, \] (A4)

\[ \text{FT}_1 = -\int_{-\pi/2}^{\pi/2} \int_{0}^{2\pi} R^2 (\sin\phi \cos\lambda \tau_{\phi} - \sin\lambda \tau_{\phi}) \cos\phi \, d\phi \, d\lambda \quad \text{and,} \] (A5)

\[ \text{FT}_2 = -\int_{-\pi/2}^{\pi/2} \int_{0}^{2\pi} R^2 (\sin\phi \sin\lambda \tau_{\phi} + \cos\lambda \tau_{\phi}) \cos\phi \, d\phi \, d\lambda, \] (A6)

where \( \text{BT}_1, \text{BT}_2, \text{MT}_1, \text{MT}_2, \text{FT}_1, \) and \( \text{FT}_2 \) are the 0° and 90°E components of the equatorial bulge, mountain, and friction torques, respectively. The quantity \( p_s \) is the surface pressure, \( R \) the earth’s radius, \( g \) the gravitational acceleration, \( \Omega \) the earth’s angular velocity, and \( h \) the topographic height. The zonal and me-
ridional components of the frictional stresses are $\tau_\lambda$ and $\tau_{dp}$ respectively. Furthermore, the two components of the EAAM vector can be written as follows (Bell 1994):

$$EM_1 = -\frac{\Omega R^4}{g} \int_0^{2\pi} \int_{-\pi/2}^{\pi/2} p_x \sin \phi \cos^2 \phi \cos \lambda \, d\phi \, d\lambda - \frac{R^3}{g} \int_0^{2\pi} \int_{-\pi/2}^{\pi/2} (u \sin \phi \cos \phi \cos \lambda - v \cos \phi \sin \lambda) \, d\phi \, d\lambda \, dp,$$

and

$$EM_2 = -\frac{\Omega R^4}{g} \int_0^{2\pi} \int_{-\pi/2}^{\pi/2} p_x \sin \phi \cos^2 \phi \sin \lambda \, d\phi \, d\lambda - \frac{R^3}{g} \int_0^{2\pi} \int_{-\pi/2}^{\pi/2} (u \sin \phi \cos \phi \sin \lambda - v \cos \phi \cos \lambda) \, d\phi \, d\lambda \, dp,$$  

where $p$ is pressure and $u$ and $v$ are the zonal and meridional winds, respectively.

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