Supernumerary Spacing of Rainbows Produced by an Elliptical-Cross-Section Cylinder. I. Theory

James A. Lock
Cleveland State University, j.lock@csuohio.edu

Follow this and additional works at: https://engagedscholarship.csuohio.edu/sciphysics_facpub

Part of the Physics Commons

How does access to this work benefit you? Let us know!

Publisher's Statement
This paper was published in Applied Optics and is made available as an electronic reprint with the permission of OSA. The paper can be found at the following URL on the OSA website: http://www.opticsinfobase.org/ao/abstract.cfm?URI=ao-39-27-5040. Systematic or multiple reproduction or distribution to multiple locations via electronic or other means is prohibited and is subject to penalties under law.

Original Citation

Repository Citation
https://engagedscholarship.csuohio.edu/sciphysics_facpub/116

This Article is brought to you for free and open access by the Physics Department at EngagedScholarship@CSU. It has been accepted for inclusion in Physics Faculty Publications by an authorized administrator of EngagedScholarship@CSU. For more information, please contact library.es@csuohio.edu.
Supernumerary spacing of rainbows produced by an elliptical-cross-section cylinder. I. Theory

James A. Lock

A sequence of rainbows is produced in light scattering by a particle of high symmetry in the short-wavelength limit, and a supernumerary interference pattern occurs to one side of each rainbow. Using both a ray-tracing procedure and the Debye-series decomposition of first-order perturbation wave theory, I examine the spacing of the supernumerary maxima and minima as a function of the cylinder rotation angle when an elliptical-cross-section cylinder is normally illuminated by a plane wave. I find that the supernumerary spacing depends sensitively on the cylinder-cross-section shape, and the spacing varies sinusoidally as a function of the cylinder rotation angle for small cylinder ellipticity. I also find that relatively large uncertainties in the supernumerary spacing affect the rainbow angle only minimally.

© 2000 Optical Society of America

OCIS codes: 290.0290, 290.4020, 080.1510.

1. Introduction

When an electromagnetic plane wave is scattered by a high-symmetry object such as a sphere, spheroid, circular-cross-section cylinder, or an elliptical-cross-section cylinder whose size is much larger than the wavelength of light, a sequence of rainbows is observed in the scattering far zone. The various rainbows are labeled by index p, with the one-internal-reflection (i.e., primary) rainbow corresponding to p = 2, the two-internal-reflection (i.e., secondary) rainbow corresponding to p = 3, etc. The rainbow is an example of the fold caustic, the simplest of the structurally stable optical caustics. A consequence of the rainbow’s structural stability is that, if a spherical particle were to be slightly deformed into a prolate or oblate spheroid, or if a cylinder’s circular cross section were to be slightly deformed into an ellipse, the resulting rainbow would distort in a number of ways. The rainbow scattering angle would shift, the spacing of the maxima and minima of the supernumerary interference pattern adjacent to the rainbow would change, and the relative intensity of the supernumerary maxima would change. But structural stability requires that the basic rainbow morphology persists.

When a circular cylinder is deformed so that its cross section becomes slightly elliptical, one of the distortions of the p = 2 and p = 3 rainbows is described by the Möbius extension to ray theory. When the distorted cylinder is rotated about its axis so that the incident beam illuminates different portions of its elliptical surface, the rainbow scattering angle, to first order in the eccentricity of the ellipse, oscillates sinusoidally back and forth about the Descartes rainbow angle for a circular cylinder.

In this paper I examine another distortion of the p = 2 and p = 3 rainbows when the cross section of a cylinder is deformed from a circle into an ellipse. The angular spacing of the maxima and minima of the supernumerary interference pattern is standardly parameterized by the quantity $h_p$, which is defined in Eq. (2) below. One observes that as the cylinder is rotated about its axis and the rainbow angle oscillates back and forth, the supernumerary pattern adjacent to it also oscillates back and forth, but by an alternately slightly larger or slightly smaller amount. This corresponds to an expansion and contraction of the supernumeraries of the p rainbow, which is described mathematically by $h_p$, being a function of the cylinder rotation angle $\xi$.

This paper proceeds as follows. In Section I briefly review the ray, Möbius, Airy, and complex angular momentum (CAM) theories of the rainbow and describe the appearance of rainbows in exact Rayleigh–Debye theory for scattering by a circular
cylinder. In Sections 3 and 4 I calculate \( h_p \) using two different methods, one in the scattering near zone and the other in the far zone. In Section 3 I use a numerical ray-tracing procedure to determine the shape of the phase fronts exiting a circular- or elliptical-cross-section cylinder in the vicinity of the \( p = 2 \) and \( p = 3 \) rainbows. From the shape of the phase fronts, I obtain \( h_{2}(\xi) \) and \( h_{3}(\xi) \) and find that the coefficients of the Fourier-series decomposition of \( h_{2}(\xi) \) and \( h_{3}(\xi) \) scale as various powers of the eccentricity of the elliptical cross section. In Section 4 I determine \( h_{3}(\xi) \) using the far-zone scattered intensity in the vicinity of the rainbow. To do this I first derive the wave theory scattering equations to first order in the perturbation of the shape of the cylinder’s cross section from that of a circle. I then perform a Debye-series decomposition of the resulting partial-wave scattering amplitudes and find that the surface shape perturbation induces a coupling between partial waves at all the interactions of the partial waves with the cylinder surface. I then numerically compute \( h_{3}(\xi) \) using first-order perturbation theory along with Airy or CAM modeling of the supernumerary intensity minima and compare the results with the ray theory results of Section 3.

Finally in Section 5 I summarize the results and comment on their significance.

2. Theories of the Rainbow

A. Rainbow in Ray Theory and Möbius Theory

I consider a family of parallel rays of wavelength \( \lambda \) and wave number \( k = 2\pi/\lambda \) normally incident on a homogeneous circular cylinder of refractive index \( n \) and radius \( a \). The scattering angle \( \theta^D_p \) of the \( p \) rainbow is given in ray theory by\(^6\)

\[
\begin{align*}
\cos(\phi^D_i) &= [(n^2 - 1)/(p^2 - 1)]^{1/2}, \\
\sin(\phi^D_i) &= (1/n)\sin(\phi^D_p), \\
\theta^D_p &= (p - 1)\pi + 2\phi^D_i - 2p\phi^D_p.
\end{align*}
\]

(1)

The same formulas are applicable to ray scattering by a sphere. The rainbow ray for a circular-cross-section cylinder or sphere is known as the Descartes ray and is denoted in Eqs. (1) by the superscript \( D \). The angle that the incident and refracted Descartes ray makes with the normal to the cylinder surface as the ray enters and exits it is \( \phi^D_i \) and \( \phi^D_p \), respectively. The shape of the phase fronts of the rays exiting a circular cylinder in the immediate vicinity of the \( p \) rainbow Descartes rays is\(^5\)

\[
Y = -h_p X^3/3a^2 + O(X^4/a^3),
\]

(2)

where \( X \) is the distance measured along the exit plane of the cylinder, defined as being tangent to the cylinder and normal to the Descartes ray; \( Y \) is the distance measured parallel to the Descartes ray; and\(^7\)

\[
h_p = [(p^2 - 1)^2(p^2 - n^2)^{1/2}]/[p^2(n^2 - 1)^{3/2}].
\]

(3)

We next consider a cylinder with an elliptical cross section whose surface is given by

\[
(x'/a')^2 + (y'/b')^2 = 1,
\]

(4)

where \( x', y', z' \) is a coordinate system attached to the cylinder whose symmetry axis coincides with the \( z' \) axis. The cylinder’s eccentricity is

\[
\epsilon = (b/a) - 1.
\]

(5)

Another set of coordinates \( x, y, z \) is fixed in the laboratory with \( z = z' \). The cylinder is oriented so that the \( x' \) axis makes an angle \( \xi \) with the \( x \) axis (this is illustrated in Fig. 1 of Ref. 8). A family of incident rays propagates in the \(-y\) direction before encountering the cylinder, and the scattering angle \( \theta \) is measured clockwise from the \(-y\) axis. The scattering angle of the \( p = 2 \) and \( p = 3 \) rainbows of the elliptical-cross-section cylinder is found in Möbius theory to be\(^4\)

\[
\begin{align*}
\theta_2^R(\xi) &= \theta_2^D - 8\epsilon\sin(\phi^D_i)\cos^3(\phi^D_p) \\
&\times \cos(2\xi + \theta_2^d) + O(\epsilon^4), \\
\theta_3^R(\xi) &= \theta_3^D + 32\epsilon\sin(\phi^D_i)\cos^3(\phi^D_p) \\
&\times \cos(2\phi^D_i)\cos(2\xi + \theta_3^d) + O(\epsilon^4). \quad (6)
\end{align*}
\]

The details of the Möbius calculation are described more fully in Ref. 8.

B. Rainbow in Airy Theory

In the physical optics model of the rainbow, also known as Airy theory, the cubic phase front of the electric field exiting a circular cylinder in the vicinity of the rainbow is Fourier transformed to the scattering far zone.\(^9\) The square of the far-zone scattered electric field is

\[
I(\theta) = (2\pi I_0 F/r)(a \lambda^{1/3}/h_p^{2/3}) \text{Ai}^2(-x^{2/3}\Delta/h_p^{1/3}),
\]

(7)

where \( I_0 \) is the intensity of the plane wave, \( x = ka \) is the cylinder size parameter, \( r \) is the distance from the cylinder axis to the far-zone position of the detector in the \( x, y \) plane, \( \text{Ai} \) is the Airy integral,\(^10\)

\[
\Delta = \theta - \theta^D_p,
\]

(8)

and \( F \) is the appropriate combination of flat-surface transmission and reflection Fresnel coefficients evaluated at the incident and transmitted angles \( \phi^D_i \) and \( \phi^D_p \) of the Descartes ray. To derive Eq. (7) it is assumed that the incident electric field is polarized parallel to the cylinder axis (i.e., the TE polarization). If the incident electric field were instead polarized perpendicular to the cylinder axis (i.e., the TM polarization), a large, if not dominant, contribution to the scattered intensity is proportional to the square of the derivative of the Airy integral, \( \text{Ai}^2(-x^{2/3}\Delta/h_p^{1/3}) \), because the internal reflections occur near the Brewster angle.\(^11\)
C. Rainbow in Rayleigh–Debye Theory and Complex Angular Momentum Theory

The exact solution for scattering of a normally incident electromagnetic plane wave by an infinitely long homogeneous circular cylinder is written in terms of an infinite series of cylindrical multipole partial waves and is sometimes called Rayleigh theory in analogy to the term Mie theory for scattering by a sphere. The decomposition of the partial-wave scattering amplitudes into the individual contributions of diffraction plus reflection \((p = 0)\), transmission \((p = 1)\), and transmission following \(p - 1 \geq 1\) internal reflections is known as the Debye series. A number of rainbows are observed in the computed Rayleigh–Debye far-zone scattered intensity for a circular-cross-section cylinder with \(x >> 1\) (e.g., the \(p = 2\) rainbow for \(n = 1.333\) and \(x = 1000.0\) is shown as the solid curve in Fig. 1). The scattered intensity in the vicinity of the rainbow, however, is modulated by scattered light because of other Debye-series contributions, such as reflection. Thus, for a careful examination of the Debye series contribution to the term Mie theory for scattering by a sphere, the Airy theory intensity of Eq. 7 for the rainbow angle \(\theta_p\). I did not pursue this calculation, however, for the following reason. The Möbius theory derivation of the rainbow angle locates the position of the Descartes ray in the exit plane of the cylinder to first order in the eccentricity \(\epsilon\). Können has extended Möbius theory for the \(p = 3\) rainbow to locate the position of the Descartes ray in the exit plane to order \(\epsilon^2\). But the determination of the shape of the exiting cubic phase front requires that the position of the Descartes ray in the exit plane be known to at least order \(\epsilon^3\).

Instead, I determined \(h_p(\xi)\) and \(h_3(\xi)\) using the following numerical ray-tracing procedure. For each cylinder rotation angle \(\xi\), I considered a family of parallel rays normally incident on the cylinder at intervals of \(\Delta \theta = 0.0002^\circ\). Each ray was propagated through the cylinder by use of the equations derived in Ref. 8, and the minimum deflection ray (i.e., the rainbow ray) was identified. Next, another family of seven rays centered on the rainbow ray at intervals of \(\Delta \theta = 0.05^\circ\) was propagated through the cylinder and the length of the optical path \(L\) of each ray from the cylinder’s entrance plane (defined as being normal to the incoming rainbow ray and tangent to the ellipse) to its exit plane (defined as being normal to the outgoing rainbow ray and tangent to the ellipse) was computed. The length of the optical

\[
E(\theta) \propto \text{Ai}[-x^{2/3} \Delta / h_p^{1/3} u(\Delta)] - iv(\Delta) x^{-1/3} \text{Ai}[-x^{2/3} \Delta / h_p^{1/3} u(\Delta)],
\]

with

\[
u(\Delta) = v_0 + v_1 \Delta + v_2 \Delta^2 + \ldots,
\]

\[
u(\Delta) = u_0 + u_1 \Delta + u_2 \Delta^2 + \ldots.
\]
path of the four rays for which $\phi_i \geq \phi^{R}_i$ was then fitted to the form

$$\frac{L}{a} = L(R)/a + A_+(X/a)^3 - h_1(X/a)^3/3 + B_+(X/a)^4,$$

(11)

where $X$ is the distance along the exit plane from the rainbow ray to the ray in question; $L(R)$ is the optical path length of the rainbow ray; and $A_+$, $h_+$, and $B_+$ are constants determined by the fitting procedure. Similarly, the length of the optical path of the four rays for which $\phi_i \leq \phi^{R}_i$ was fitted to the form

$$\frac{L}{a} = L(R)/a + A_-(X/a)^3 - h_-(X/a)^3/3 + B_-(X/a)^4.$$

(12)

Because the rainbow ray was located to within $\Delta \phi_i \leq 0.0002^\circ$, the values of $A_+$ and $A_-$ were found to be exceedingly close to zero as required by Eq. (2). Because the $\Delta \phi_i$ interval for the set of seven rays near the rainbow ray was small, a correspondingly small portion of the exiting phase front was sampled. As a result, $h_+$ and $h_-$ were found to be nearly identical, and $B_+$ and $B_-$ differed only minimally. The supernumerary scaling parameter $h_p$ was then taken to be

$$h_p(\xi) = (h_+ + h_-)/2.$$

(13)

After this procedure was carried out for $0^\circ \leq \xi \leq 360^\circ$ in intervals of $\Delta \xi = 1^\circ$ for both the $p = 2$ and $p = 3$ rainbows, the resulting functions $h_2(\xi)$ and $h_3(\xi)$ were decomposed into the Fourier series

$$h_2(\xi) = e_0 + \sum_{m=1}^{\infty} e_m \cos(m\xi) + \sum_{m=1}^{\infty} f_m \sin(m\xi),$$

(14)

$$h_3(\xi) = g_0 + \sum_{m=1}^{\infty} g_m \cos(m\xi) + \sum_{m=1}^{\infty} j_m \sin(m\xi).$$

(15)

An analogous Fourier-series decomposition of the rainbow angles $\theta_2^{R}(\xi)$ and $\theta_3^{R}(\xi)$ was performed in Ref. 8. The Fourier decomposition of $h_2(\xi)$ and $h_3(\xi)$ provides a useful diagnostic for our choice of $\Delta \phi_i$ for the family of incident rays used to locate the rainbow ray and for the family of seven rays used to determine the shape of the exiting phase front. If the cylinder has a circular cross section, all the Fourier coefficients for $m \neq 0$ should vanish. Similarly, if the cylinder has an elliptical cross section, all the odd-$m$ Fourier coefficients should vanish because of the $180^\circ$ rotational symmetry of the ellipse. I found that if $\Delta \phi_i$ were much larger than 0.0002$^\circ$ for the initial family of incident rays, the rainbow ray was not located with sufficient precision in the exit plane. As a result, the fitted values of $A_+$ and $A_-$ were no longer near zero, and the odd-$m$ Fourier coefficients exhibited a substantial amount of noise. Similarly, if $\Delta \phi_i$ was much larger than 0.05$^\circ$ for the family of seven rays in the vicinity of the rainbow ray, a larger than desired portion of the exiting wave front was sampled. The fitting of this larger portion to only order $(X/a)^3$ did not allow $h_2(\xi)$ and $h_3(\xi)$ to be determined with sufficient precision because higher powers of $X/a$ also contribute substantially to the shape of the phase front over the larger interval. As a consequence, the odd-$m$ Fourier coefficients again exhibited a substantial amount of noise.

The Fourier coefficients of $h_2(\xi)$ and $h_3(\xi)$ exhibited the following scaling behavior in $\epsilon$. For a circular-cross-section cylinder, $e_0$ and $g_0$ are equal to $h_2$ and $h_3$, respectively. For an elliptical-cross-section cylinder, both $e_0$ and $g_0$ differ from the circular-cross-section value of $h_2$ and $h_3$ of Eq. (3) by a term approximately linear in $\epsilon$, and the even-$m$ coefficients $e_m, f_m, g_m$, and $j_m$ for $m \geq 2$ are approximately proportional to $\epsilon^{m/2}$, with the constants of proportionality depending on refractive index. An identical scaling behavior in $\epsilon$ was previously found for the Fourier coefficients of the rainbow angle $\theta_2^{R}(\xi)$ and $\theta_3^{R}(\xi)$. As to refractive-index dependence, the $m = 2$ portion of Eq. (14) for the $p = 2$ rainbow was found to be well fit by the expression

$$h_2(\xi) \approx e_0 + 19\epsilon [\sin(\phi_i^D)]^{1/4} \cos(\phi_i^D) \cos(2\xi + \Phi),$$

(16)

where the difference between $e_0$ and $h_2$ of Eq. (3) is approximately an order of magnitude smaller than the amplitude of the $\cos(2\xi + \Phi)$ term in approximation (16), and

$$\Phi = (250^\circ)n - 285^\circ.$$  

(17)

For $1.25 \leq n \leq 1.7$, the amplitude of the $\cos(2\xi + \Phi)$ term agrees with the results of the numerical ray-tracing calculation to better than 2%, and $\Phi$ agrees to within approximately 5°. I attribute no fundamental theoretical basis to approximations (16) and (17). They are solely the result of our motivation to obtain an approximate equation for $h_2(\xi)$ to first order in $\epsilon$ that resembles the Möbius theory result of Eqs. (6) for $\theta_2^{R}(\xi)$.

In Fig. 2 I plot $h_2(\xi) - h_{\text{ave}}$ as a function of $\xi$ for an elliptical-cross-section cylinder with $\epsilon = 0.0001$ and $n = 1.333$, where $h_{\text{ave}}$ is the average value of $h_2(\xi)$ over the interval $0^\circ \leq \xi \leq 180^\circ$. The behavior of $h_2(\xi) - h_{\text{ave}}$ for $180^\circ \leq \xi \leq 360^\circ$ is identical to that for $0^\circ \leq \xi \leq 180^\circ$ because of the $180^\circ$ rotational symmetry of the elliptical cross section. The small jaggedness of the solid curve is due to the approximate nature of our numerical ray-tracing procedure. The dashed curve is the $m = 2$ Fourier component of $h_2(\xi) - h_{\text{ave}}$ that, according to the scaling behavior in $\epsilon$ mentioned in the above paragraph, is the dominant contribution in the Möbius regime, $\epsilon \ll 1$. The results for $h_3(\xi)$ were similar to those of Fig. 2 and, together with Fig. 2, illustrate that to determine the shape of the exiting phase front, and from it $h_3(\xi)$, by numerical ray tracing requires the rainbow ray to be located accurately in the exit plane and the exiting phase front to be sampled over a sufficiently small interval about the rainbow ray.
4. Supernumerary Spacing Parameter \( h_p \) in First-Order Perturbation Rayleigh–Debye Theory

The computation of light scattering by a nonspherical particle or by a cylinder with a noncircular cross section is best handled for moderate size parameters and for moderate eccentricities by the so-called T-matrix method.\(^{21,22}\) This approach, however, is not appropriate for a detailed quantitative study of rainbows because the cylinder size parameter that is required, \( x \sim 1000 \) or greater, far exceeds the size parameter range for which T-matrix computations are numerically stable.\(^{23}\) As an alternative, we employ the first-order perturbation theory extension to the exact Rayleigh–Debye partial-wave scattering equations. In Refs. 24–27, the equations for wave scattering by a nonspherical particle were derived to first order in the particle-shape perturbation, and in Ref. 28 the form of the equations was simplified substantially. Similarly, in Ref. 29 the equations for wave scattering by a cylinder with a noncircular cross section were derived to first order in the cylinder-shape perturbation. These first-order perturbation-theory equations can again be simplified as follows.

The surface of the cylinder is taken to be

\[
 r(\theta) = a[1 + \delta f(\theta)]
\]

(18)

in polar coordinates where \( a \delta f(\theta) \) is the surface-shape perturbation from a circle, normalized so that

\[
f(\theta) = O(1).
\]

(19)

First-order perturbation theory is an accurate approximation to exact wave scattering theory when\(^{30}\)

\[
 \delta \ll 1,
\]

(20)

\[
 ka \delta \ll 2\pi.
\]

(21)

Inequality (21) is usually the more restrictive of the two and limits the regime of applicability of first-order perturbation theory for large-radius cylinders to small perturbations.

The geometry of the plane wave used here is different from that for the ray theory in Section 3. The electric field strength of the plane wave is \( E_0 \) and it is normally incident on the cylinder with its propagation direction parallel to the \( x \) axis and its electric field polarized either parallel to the \( z \) axis (TE polarization) or parallel to the \( y \) axis (TM polarization). We define the polarization-dependent quantities \( \alpha \) and \( \beta \) by

\[
 \alpha = \begin{cases} 
 n & \text{for TE}, \\
 1 & \text{for TM}, 
\end{cases}
\]

(22)

\[
 \beta = \begin{cases} 
 1 & \text{for TE} \\
 n & \text{for TM} 
\end{cases}
\]

and let

\[
y = nka,
\]

(23)

\[
 E_i = \alpha H^{(1)}_l(x)J_l(y) - \beta H^{(1)}_l(x)J_l(y),
\]

(24)

where \( J_l \) are Bessel functions and \( H^{(1)}_l \) and \( H^{(2)}_l \) are Hankel functions of the first and second kind. The time dependence of the plane wave is \( \exp(-i\omega t) \) so that \( H^{(1)}_l \) and \( H^{(2)}_l \) describe radially outgoing and incoming cylindrical multipole waves, respectively.\(^{31}\) By matching the tangential components of the incident, scattered, and interior electric and magnetic fields \( E_x, B_y, E_y + \delta(dE/d\theta) E_x, \) and \( B_x + \delta(dB/d\theta) B_y \) at the surface of the cylinder to first order in \( \delta \), we obtain the far-zone scattered intensity:

\[
 I(\theta) = (2I_0/\pi kr)|F(\theta)|^2,
\]

(25)

where the scattering amplitude \( F(\theta) \) is

\[
 F(\theta) = \sum_{l = -\infty}^{\infty} b_l \exp(i\ell\theta)
\]

(26)

for a TE incident plane wave and

\[
 F(\theta) = \sum_{l = -\infty}^{\infty} a_l \exp(i\ell\theta)
\]

(27)

for a TM incident plane wave. The partial-wave scattering amplitudes \( a_l \) and \( b_l \) for a circular cylinder are given by

\[
 \begin{bmatrix} a_l^{(0)} \\ b_l^{(0)} \end{bmatrix} = \begin{bmatrix} \alpha J_l(x)J_l(y) - \beta J_l(x)J_l(y) \end{bmatrix}/E_t,
\]

(28)

where the superscript \((0)\) henceforth indicates a circular cross section. For completeness, the TE and TM partial-wave interior amplitudes are

\[
 \begin{bmatrix} c_l^{(0)} \\ d_l^{(0)} \end{bmatrix} = 2i/(n^2 \pi x E_l).
\]

(29)
When the surface of the cylinder is noncircular, the partial wave-scattering and interior amplitudes become, after a large amount of algebra,

\[ a_t = a_t^{(0)} + [2i(n^2 - 1)\delta/\pi] \]
\[ \times \sum_{l=-\infty}^{\infty} i^{l-l} a_t^{(1)} + O(\delta^2), \]
\[ b_t = b_t^{(0)} + [2i(n^2 - 1)\delta/\pi] \]
\[ \times \sum_{l=-\infty}^{\infty} i^{l-l} b_t^{(1)} + O(\delta^2), \]
\[ c_t = c_t^{(0)} - [2i(n^2 - 1)\delta/(n^2\pi)] \]
\[ \times \sum_{l=-\infty}^{\infty} i^{l-l} c_t^{(1)} + O(\delta^2), \]
\[ d_t = d_t^{(0)} - [2i(n^2 - 1)\delta/(n^2\pi)] \]
\[ \times \sum_{l=-\infty}^{\infty} i^{l-l} d_t^{(1)} + O(\delta^2), \]  \tag{30}

where the superscript (1) henceforth indicates a first-order correction in the perturbation strength $\delta$. The first-order corrections are

\[ a_{t,i}^{(1)} = \{[(l'/\chi^2)J_{l'}(y)J_{l}(y) \]
\[ + J_{l'}(y)J_{l}(y)]I_{l,i} \]
\[ - (il'/\chi^2)J_{l'}(y)J_{l}(y)I_{l',i}]/(E_lE_i), \]
\[ b_{t,i}^{(1)} = [J_{l'}(y)J_{l}(y)]I_{l,i}/(E_lE_i), \]
\[ c_{t,i}^{(1)} = \{[(l'/\chi^2)J_{l'}(y)H_{l'}^{(1)}(x) \]
\[ + J_{l'}(y)H_{l'}^{(1)}(x)]I_{l',i} - (il'/\chi^2)J_{l'}(y) \]
\[ \times H_{l'}^{(1)}(x)I_{l,i}]/(E_lE_i), \]
\[ d_{t,i}^{(1)} = [J_{l'}(y)H_{l'}^{(1)}(x)]I_{l,i}/(E_lE_i), \]  \tag{31}

where

\[ I_{l,i} = (1/2\pi) \int_{0}^{2\pi} d\theta f(\theta) \exp[i(l' - l)\theta], \]
\[ I_{l,i}' = (1/2\pi) \int_{0}^{2\pi} d\theta (df/d\theta) \exp[i(l' - l)\theta]. \]  \tag{32}

The Fourier-series decomposition of the surface perturbation $f(\theta)$ is

\[ f(\theta) = A_0 + \sum_{q=1}^{\infty} A_q \cos[q(\theta - \xi)] \]
\[ + \sum_{q=1}^{\infty} B_q \sin[q(\theta - \xi)], \]  \tag{33}

where again $\xi$ is the rotation angle of the cylinder’s $x'$ axis with respect to the laboratory $x$ axis. Inserting Eq. (33) into Eqs. (32), we obtain

\[ I_{l,i} = A_q \delta_{l',i} + (1/2) \sum_{q=1}^{\infty} (A_q - iB_q) \exp(-iq\xi) \delta_{l',i}, \]
\[ + (1/2) \sum_{q=1}^{\infty} (A_q + iB_q) \exp(iq\xi) \delta_{l',i+q}, \]
\[ I_{l,i}' = (1/2) \sum_{q=1}^{\infty} iq(A_q - iB_q) \exp(-iq\xi) \delta_{l',i}, \]
\[ - (1/2) \sum_{q=1}^{\infty} iq(A_q + iB_q) \exp(iq\xi) \delta_{l',i+q}, \]  \tag{34}

where $\delta_{l',i}$ is the Kronecker delta symbol. Equations (34) illustrate that two different incident partial waves, $l' = l + q$ and $l' = l - q$, are coupled to each scattered partial wave $l$ by the Fourier component $q$ of the surface perturbation. This is a simpler situation than for scattering by a perturbed spheroid where all the incident partial waves $l'$ in the interval $l - q \leq l' \leq l + q$ are coupled with various strengths to each scattered partial wave $l$ by the Fourier component $q$ of the surface perturbation.

Before we can apply first-order perturbation theory to the $p$ rainbow of an elliptical-cross-section cylinder illuminated by a TE-polarized plane wave, we must first perform a Debye-series decomposition of the partial-wave scattering amplitudes and retain only the $p$ term in Eq. (26) so as to rid the scattered intensity of the influence of other scattering processes such as reflection. We denote the cylinder interior as region 1 and the exterior as region 2. We consider a radially incoming TE or TM cylindrical multipole wave in region 2 with the partial-wave number $l'$. When it encounters the cylinder surface, a portion $\mathcal{R}_{l',l}^{21}$ of the wave amplitude is transmitted into the cylinder as a different cylindrical multipole wave with the partial-wave number $l$. The remaining portion of the wave amplitude $\mathcal{R}_{l',l}^{22}$ is reflected from the surface as a radially outgoing cylindrical multipole wave with the partial-wave number $l$. Following the procedure in the appendix of Ref. 20 applied to the perturbed cylinder geometry, we obtain

\[ \mathcal{R}_{l',l}^{21} = T_{l',l} \delta_{l',l} - [4i(n^2 - 1)\delta/\pi]l' - l \mathcal{T}_{l,l}^{21} + O(\delta^2), \]
\[ \mathcal{R}_{l',l}^{22} = R_{l',l} \delta_{l',l} - [4i(n^2 - 1)\delta/\pi]l' - l \mathcal{R}_{l,l}^{22} + O(\delta^2), \]  \tag{35}

where

\[ T_{l}^{21} = 4i/(n\pi D_l), \]
\[ R_{l}^{22} = -[\alpha H_{l}^{22}(x)H_{l}^{22}(y) - \beta H_{l}^{22}(x)H_{l}^{22}(y)]/D_l, \]  \tag{36}

with

\[ D_l = \alpha H_{l}^{21}(x)H_{l}^{21}(y) - \beta H_{l}^{21}(x)H_{l}^{21}(y). \]  \tag{37}

Equations (36) and (37) describe the conservation of partial-wave number for transmission and reflection.
of a cylindrical multipole wave by a circular-cross-section cylinder and correspond to conservation of angular momentum in analogous expressions for quantum mechanical scattering. In addition, the quantities

\[
T_{r,l}^{21} = \left\{ [(l^2/xy)H_{l}^{(2)}(y)H_{l}^{(1)}(x) + H_{l}^{(2)}(y)H_{l}^{(1)}(x)]I_{r,l} - (il'/xy) \times H_{l}^{(2)}(y)H_{l}^{(1)}(x)I_{r,l} \right\}/(nD)_{D},
\]

\[
R_{r,l}^{22} = \left\{ [(l^2/xy)H_{l}^{(2)}(y)H_{l}^{(2)}(y) + H_{l}^{(2)}(y)H_{l}^{(2)}(y)]I_{r,l} - (il'/xy)H_{l}^{(2)}(y) \times H_{l}^{(2)}(y)I_{r,l} \right\}/(D)_{D} \tag{38}
\]

for the TM polarization and

\[
T_{r,l}^{21} = [H_{l}^{(2)}(y)H_{l}^{(1)}(x)]I_{r,l}/(nD)_{D},
\]

\[
R_{r,l}^{22} = [H_{l}^{(2)}(y)H_{l}^{(2)}(y)]I_{r,l}/(D)_{D} \tag{39}
\]

for the TE polarization describe the coupling between the partial waves \(l'\) and \(l\) of the incident and scattered light induced by the noncircular character of the surface shape. In quantum mechanical scattering, analogous expressions describe the target taking up some of the angular momentum of the projectile during the scattering process.

Similarly, when a TE or TM radially outgoing cylindrical multipole wave in region 1 with the partial-wave number \(l'\) encounters the cylinder surface, a portion \(\mathcal{R}_{r,l}^{21}\) of the wave amplitude is transmitted to the exterior as a different cylindrical multipole wave with the partial-wave number \(l\), and the remaining portion of the wave amplitude \(\mathcal{R}_{r,l}^{11}\) is reflected back inside the cylinder as a different cylindrical multipole wave with the partial-wave number \(l\). Again, following the procedure in the appendix of Ref. 20 applied to the perturbed cylinder geometry, we obtain

\[
\mathcal{R}_{r,l}^{12} = T_{r,l}^{12} \mathcal{R}_{r,l}^{11} - [4i(n^2 - 1)\delta/\pi]T_{r,l}^{12} + O(\delta^2),
\]

\[
\mathcal{R}_{r,l}^{11} = R_{r,l}^{11} \mathcal{R}_{r,l}^{11} - [4i(n^2 - 1)\delta/\pi]R_{r,l}^{11} + O(\delta^2), \tag{40}
\]

where

\[
T_{r,l}^{12} = 4in/(\pi xD),
\]

\[
R_{r,l}^{11} = -[\alpha H_{l}^{(1)}(x)H_{l}^{(1)}(x) - \beta H_{l}^{(1)}(x)H_{l}^{(1)}(x)]/D, \tag{41}
\]

\[
T_{r,l}^{12} = \left\{ [(l^2/xy)H_{l}^{(1)}(x)H_{l}^{(1)}(y) + H_{l}^{(1)}(x)H_{l}^{(1)}(y)]I_{r,l} - (il'/xy)H_{l}^{(1)}(x) \times H_{l}^{(1)}(y)I_{r,l} \right\}/[n/(D)D],
\]

\[
R_{r,l}^{11} = \left\{ [(l^2/xy)H_{l}^{(1)}(x)H_{l}^{(1)}(y) + H_{l}^{(1)}(x)H_{l}^{(1)}(y)]I_{r,l} - (il'/xy)H_{l}^{(1)}(x) \times H_{l}^{(1)}(y)I_{r,l} \right\}/(D)D \tag{42}
\]

for the TM polarization and

\[
T_{r,l}^{12} = [H_{l}^{(1)}(x)H_{l}^{(1)}(y)]I_{r,l}/[n/(D)D],
\]

\[
R_{r,l}^{11} = [H_{l}^{(1)}(x)H_{l}^{(1)}(y)]I_{r,l}/(D)D \tag{43}
\]

for the TE polarization.

Now that we have obtained the partial-wave transmission and reflection Fresnel coefficients to first order in \(\delta\), again after a great amount of algebra, one finds that the partial-wave scattering amplitudes of Eqs. (28) and (31) can be written in terms of these Fresnel coefficients as

\[
\mathbf{a}_{l}^{(0)} \mathbf{b}_{l}^{(0)} = \frac{1}{2} \left[ 1 + R_{l}^{22} - T_{l}^{21}(1 - R_{l}^{11})^{-1}T_{l}^{12} \right] = \frac{1}{2} \left[ 1 + R_{l}^{22} - \sum_{p=1}^{n} T_{l}^{21} R_{l}^{11}^{-p-1} T_{l}^{12} \right], \tag{44}
\]

\[
\mathbf{a}_{l}^{(1)} \mathbf{b}_{l}^{(1)} = R_{l}^{22} + \sum_{p=0}^{n} T_{l}^{21} R_{l}^{11}^{-p-1} T_{l}^{12} \tag{45}
\]

where the TE version of the coefficients is to be used in the decomposition of \(\mathbf{b}_{l}^{(0)}\) and \(\mathbf{b}_{l}^{(1)}\) and the TM version is to be used in the decomposition of \(\mathbf{a}_{l}^{(0)}\) and \(\mathbf{a}_{l}^{(1)}\). Equations (44) are the original series expansion of Debye,\(^{13}\) and Eqs. (45) are its extension to first order in \(\delta\) for a cylinder with a noncircular cross section. The significance of Eqs. (45) is that it clearly illustrates the physical mechanism by which the noncircular component of the surface shape induces a coupling between the incoming and outgoing partial waves \(l'\) and \(l\). The coupling is produced at any of the interactions of the cylindrical multipole waves with the surface, i.e., at the external reflection, at the initial transmission into the cylinder, at any of the internal reflections, and at the final transmission out of the cylinder. Because the partial-wave scattering amplitudes are calculated to only first order in \(\delta\), only one change of the partial-wave number is permitted in each term of Eqs. (45). The extension of perturbation theory to \(O(\delta^2)\) would include up to \(n\) changes of the partial-wave number in each term.
of the partial-wave interior amplitudes of Eqs. (29) and (31) is found after much algebra to be

\[
\frac{c_j^{(0)}}{d_j^{(0)}} = (1/n)T_l^{21}(1 - R_l^{11})^{-1} = (1/n) \sum_{p=1}^{\infty} T_l^{21}(R_l^{11})^{-p-1},
\]

(46)

\[
\frac{c_j^{(1)}}{d_j^{(1)}} = 2nT_{l,j}^{21}(1 - R_l^{11})^{-1} + 2nT_{l,j}^{21} \\
\times (1 - R_l^{11})^{-1}R_{l,j}^{11}(1 - R_l^{11})^{-1} = 2n \sum_{p=1}^{\infty} T_{l,j}^{21}(R_l^{11})^{-p-1} \\
+ 2n \sum_{s=0}^{\infty} \sum_{t=0}^{\infty} T_{l,j}^{21}(R_l^{11})R_{l,j}^{11}(R_l^{11})^t,
\]

(47)

where again the TE version of the Fresnel coefficients is to be used in the decomposition of \(d_j^{(0)}\) and \(d_j^{(1)}\), and the TM version is to be used in the decomposition of \(c_j^{(0)}\) and \(c_j^{(1)}\). The partial-wave coupling induced by the noncircular component of the surface shape is implicit in the \(T\)-matrix formalism \(^{22,33}\) and provides the physical explanation for the optimal positioning of a tightly focused laser beam on a microparticle when morphology-dependent resonances are excited \(^{28,34,35}\).

Equations (45) are not especially convenient for numerical computations because they contain Bessel functions of negative order. If we assume that the cylinder surface shape is given in Fourier-series form and we substitute Eqs. (34) into Eqs. (31), (38), and (42), the scattering amplitude \(F(\theta)\) for TE scattering can be simplified to

\[
F(\theta) = F^{(0)}(\theta) + [2i(n^2 - 1)\delta/\pi]A_q Q_\delta(\theta) \\
+ [2i(n^2 - 1)\delta/\pi] \sum_{q=1}^{\infty} i^q Q_\delta(q, \theta) \\
\times [A_q \cos(q\xi) - B_q \sin(q\xi)] \\
- [2i(n^2 - 1)\delta/\pi] \sum_{q=1}^{\infty} i^q Q_\delta(q, \theta) \\
\times [A_q \sin(q\xi) + B_q \cos(q\xi)].
\]

(48)

The first term of Eq. (48),

\[
F^{(0)}(\theta) = b_0^{(0)} + 2 \sum_{l=1}^{\infty} b_l^{(0)} \cos(l\theta),
\]

(49)

is the scattering amplitude for a circular-cross-section cylinder in Rayleigh–Debye theory, and the expressions

\[
Q_\delta(q, \theta) = U_0 + 2 \sum_{l=1}^{\infty} U_l \cos(l\theta),
\]

\[
Q_\delta(q, \theta) = \sum_{l=0}^{\infty} V_{l,q}^+\{\cos(l\theta) + (-1)^q \cos[(l + q)\theta]\} \\
+ \sum_{l=0}^{q-1} (-1)^q V_{l,q}^- \cos(l\theta),
\]

\[
Q_\delta(q, \theta) = \sum_{l=0}^{\infty} V_{l,q}^+\{\sin(l\theta) - (-1)^q \sin[(l + q)\theta]\} \\
- \sum_{l=1}^{q-1} (-1)^q V_{l,q}^- \sin(l\theta)
\]

(50)

are linear in \(\delta\) in Eq. (48) with

\[
U_l = J_l^2(y)/E_l^2, \\
V_{l,q}^+ = J_l(y)\tilde{J}_{q,l}(y)/(E_lE_{q,l}), \\
V_{l,q}^- = J_l(y)\tilde{J}_{q,l}(y)/(E_lE_{q,l})
\]

(51)

are the contribution to the scattering amplitude provided by the various Fourier components of the surface perturbation. The term proportional to \(A_q\) corresponds to a perturbation in the form of a small increase of radius, and the terms proportional to \(A_q\) and \(B_q\) for \(q \geq 1\) correspond to a shape perturbation. The one-internal-reflection portion of Eqs. (51) is

\[
U_l = T_{l,l}^{21}R_{l,l}^{11}T_{l,l}^{12} + T_{l,l}^{21}R_{q,l,l}^{11}T_{l,l}^{12} \\
+ T_{l,l}^{21}R_{q,l,l}^{11}T_{l,l}^{12}, \\
V_{l,q}^+ = T_{q,l,l}^{21}R_{q,l,l}^{11}T_{q,l,l}^{12} + T_{q,l,l}^{21}R_{q,l,l}^{11}T_{q,l,l}^{12} \\
+ T_{q,l,l}^{21}R_{q,l,l}^{11}T_{q,l,l}^{12}, \\
V_{l,q}^- = T_{q,l,l}^{21}R_{q,l,l}^{11}T_{q,l,l}^{12} + T_{q,l,l}^{21}R_{q,l,l}^{11}T_{q,l,l}^{12} \\
+ T_{l,l}^{21}R_{q,l,l}^{11}T_{q,l,l}^{12}.
\]

(52)

The first, second, and third terms of \(V_{l,q}^+\) and \(V_{l,q}^-\) correspond to the change in the partial-wave number occurring at the initial refraction into the cylinder, at the internal reflection, and at the final refraction out of the cylinder, respectively. Because \(U_l\) corresponds only to an increase of radius, the partial-wave number is conserved at the initial refraction, at the internal reflection, and at the final refraction.

We now specify the cylinder-shape Fourier coefficients \(A_q\) and \(B_q\). When the cylinder rotation angle is \(\xi = 0\), the ellipse of Eq. (4) can be expressed in polar coordinates as

\[
r(\theta) = a(1 + \epsilon)[1 + (2\epsilon + \epsilon^2)\cos^2(\theta)]^{-1/2},
\]

(53)

and the surface perturbation is

\[
f(\theta) = r(\theta) - a.
\]

(54)

The first few Fourier coefficients of Eq. (54) are

\[
A_0 = \epsilon/2 - 3\epsilon^3/16 + 3\epsilon^5/32 + \ldots, \\
A_2 = -\epsilon^2/2 + 5\epsilon^3/64 + \ldots, \\
A_4 = 3\epsilon^4/16 - 3\epsilon^6/32 + \ldots, \\
A_6 = -5\epsilon^5/64 + \ldots, \\
B_2 = B_4 = B_6 = 0.
\]

(55)

The leading term of the coefficients \(A_q\) for \(q \geq 2\) is proportional to \(\epsilon^{q/2}\). Because first-order perturbation theory is accurate to only order \(\epsilon\), the association

\[
\delta = \epsilon/2, \quad A_0 = 1, \quad A_2 = -1
\]

(56)
describes the ellipse in first-order perturbation theory in a self-consistent way. At this level of approximation, an ellipse and a quadrupole deformation of a circle are equivalent. Similarly, a perturbation in the form of a small increase of radius is

$$\delta = \epsilon/2, \quad A_0 = 1.$$  \hfill (57)

In the numerical studies of these equations, I first performed the following test as a check of my analytical perturbation-theory formulas, the computer program that implements them, and to explore the limit of applicability of first-order perturbation theory. I computed the $p = 2$ portion of the scattered intensity for a circular-cross-section cylinder having $n = 1.333$ and $x = 1000.0$ using exact Rayleigh–Debye theory. The angle of the intensity minimum between the first and second supernumerary maxima $\theta_a^{\text{min}}$ and the angle of the intensity minimum between the second and third supernumerary maxima $\theta_b^{\text{min}}$, both measured in degrees, were determined to within $\pm 0.000 005^\circ$. I then assumed that these angles correspond to the first and second zeros of the Airy integral, leading to the values of $A_i(2.338 107)$ and $A_i(-4.087 949)$, respectively. I determined the rainbow angle $\theta_a^{R}$ and the supernumerary spacing parameter $h_2$ from Eqs. (7) and (8) by

$$2.338 107 = (\pi/180)x^{2/3}(\theta_a^{\text{min}} - \theta_a^{R})/h_2^{1/3},$$
$$4.087 949 = (\pi/180)x^{2/3}(\theta_b^{\text{min}} - \theta_b^{R})/h_2^{1/3},$$  \hfill (58)

or

$$\theta_a^{R} = (2.336 182)\theta_a^{\text{min}} - (1.336 182)\theta_a^{\text{min}};$$
$$h_2 = (0.992 283)(10^{-6})x^{2/3}(\theta_b^{\text{min}} - \theta_a^{\text{min}})^3.$$  \hfill (59)

Using Eqs. (59), I found $\theta_a^{R} = 137.945 68^\circ$ and $h_2 = 4.643 221,$ which differ only slightly from the ray theory values $\theta_a^{R} = 137.921 89^\circ$ and $h_2 = 4.899 194$ of Eqs. (1) and (3). The exact wave theory results differ slightly from the ray theory results because we are not in the $x \to \infty$ limit and because Rayleigh–Debye intensity is not perfectly fit by the square of an Airy integral. In particular, curved-surface Fresnel transmission and reflection coefficients rather than flat-surface Fresnel coefficients should be used in Eq. (7), the angle dependence of the Fresnel coefficients should be included, and terms of higher order than $X^3/a^2$ in the phase-front shape of Eq. (2) should be included.

These improvements are taken into account in CAM theory; when we model the intensity minima of the $p = 2$ rainbow using the lowest-order CAM correction to Airy theory, we obtain an improved agreement between first-order perturbation theory and ray theory. In Ref. 17, Khare and Nussenzveig obtained $u_1 = 0.202$ and $v_0 = 0.473$ for electromagnetic scattering of a TE-polarized plane wave by a sphere with $n = 1.33$. We use these values of $u_1$ and $v_0$ for scattering by both a circular cylinder, which should be appropriate, and by an elliptical-cross-section cylinder, which may not be entirely appropriate. The presence of the $A_i'$ term in relation (9) is of no consequence in the determination of $\theta_a^{\text{min}}$ and $\theta_b^{\text{min}}$ because the relative maxima and minima of Eqs. (7) and the intensity corresponding to relation (9) are identical. But when the progressively increased stretching of the exact Rayleigh–Debye wave theory intensity with respect to Airy theory is taken into account by replacing $x^{2/3}/h_p^{1/3}$ in Eqs. (58) with $(x^{2/3}/h_p^{1/3})(1 + 0.202\delta)$, we obtain $\theta_a^{R} = 137.914 91^\circ$ and $h_2 = 4.954 462$. This makes up approximately 71% and 78%, respectively, of the difference between ray theory and first-order perturbation theory with Airy modeling of the supernumerary intensity minima.

I then determined the limit to the validity of first-order perturbation theory as follows. I repeated the procedure of Eqs. (59) using Airy theory to model the supernumerary intensity minima for a set of slightly larger circular cylinders using two different methods. First, I used exact Rayleigh–Debye theory for $n = 1.333$ and $x = (1000.0)(1 + \delta)$ with $\delta = 0.001$. Second, for the same set of circular cylinders I used first-order perturbation theory with $n = 1.333$, $x = 1000.0$, and with the perturbation being an increase of radius as in Eqs. (57). The values of $\theta_a^{\text{min}}$ and $\theta_b^{\text{min}}$ that I obtained using first-order perturbation theory and exact Rayleigh–Debye theory agreed to within $\pm 0.000 02^\circ$ for $\delta \leq 0.0001$, or for $1000.0 \leq x \leq 1000.1$, which is in agreement with inequality (21). As long as $\delta \leq 0.0001$, both minimum intensity angles in perturbation theory increased by the same amount with respect to their Rayleigh–Debye values. So the resulting value of $h_2$ from Eqs. (59) remained identical to the Rayleigh–Debye result. But for much larger values of $\delta$, the two minimum intensity angles in perturbation theory increased by differing amounts, leading to differences between the perturbation-theory and Rayleigh–Debye values of $h_2$. Thus I consider $\delta = 0.0001$ for $x = 1000.0$ to be a safe upper limit for the accuracy of first-order perturbation theory to determine $h_2$ for a circular-cross-section cylinder.

I also tried to determine $h_2$ from the first two supernumerary maxima of the perturbation-theory intensity. Again I associated these maxima with the first maximum and the first minimum of the Airy integral, leading to the values of $A_i(-1.018 793)$ and $A_i(-3.248 198)$, respectively. Using this method, I found that the safe upper limit for first-order perturbation theory was now an order of magnitude smaller than it was when the intensity minima were used. This is due to the fact that first-order perturbation theory introduced a slowly varying background intensity superposed on the rainbow. This shifted the angles of the broad intensity maxima of Fig. 1 by larger differing amounts than it did for the sharp intensity minima.

Thus I obtained the best performance from first-order perturbation theory when I analyzed the first two intensity minima of the supernumerary interference pattern.

I then used first-order perturbation theory with Eqs. (56) and (59) and with Airy modeling of the
intensity minima to obtain \( \theta_2^R(\xi) \) and \( h_2(\xi) \) for an elliptical-cross-section cylinder having \( n = 1.333, x = 1000.0, \epsilon = 0.0001, \) and \( \delta = 0.00005 \) as a function of the cylinder rotation angle \( \xi \). For the Airy theory modeling, we assumed that the scattered intensity in the vicinity of the rainbow is given by

\[
I(\theta, \xi) \propto \Lambda^2 \left( -x_{ave}^{2/3} \theta - \theta_p^R(\xi) \right)^2 / p^R(\xi)^{1/3},
\]

where \( x_{ave} \) is the average value of the size parameter of the elliptical cross section. The resulting behavior of \( h_2(\xi) - h_{ave} \) and \( \theta_2^R(\xi) - \theta_{ave}^R \) is shown in Figs. 2 and 3, respectively. In obtaining these graphs I took into account the difference in the geometry of the plane wave in ray theory and Rayleigh–Debye theory.

In ray theory the plane wave propagates in the \(-y\) direction, it is incident on a cylinder with eccentricity \( \epsilon \), and the scattering angle \( \theta \) is measured clockwise from the \(-y\) axis; however, in first-order perturbation Rayleigh–Debye theory the incident plane wave propagates in the \(x\) direction so that it sees a cylinder of eccentricity \( -\epsilon \), and the scattering angle is measured counterclockwise from the \(x\) axis so that the observer sees the angle \( \theta \). In Fig. 3, \( \theta_p^R(\xi) - \theta_{ave}^R \) is compared with that obtained with the ray-tracing procedure of Section 2. The results are virtually identical and agree exactly with the Möbius theory prediction of Eqs. (6). The value of \( \theta_{ave}^R \) for the elliptical-cross-section data of Fig. 3 is 137.945 39° which compares favorably with \( \theta_p^R = 137.945 68° \) obtained from exact Rayleigh–Debye theory for a circular-cross-section cylinder with \( x_{ave} = 1000.05 \).

In Fig. 2 for \( h_2(\xi) - h_{ave} \) the prediction of ray theory is the solid curve, the \( m = 2 \) Fourier component of the ray theory result is the dashed curve, and the results of first-order perturbation theory with Airy theory modeling of the rainbow intensity minima are the open circles. The comparison between the first-order perturbation theory and ray theory results is not as close as it was for the rainbow angle. The amplitude of the oscillation of \( h_2 \) in perturbation theory is approximately 17% higher than that for ray theory, the oscillation is approximately 20° out of phase with ray theory, and the oscillation is not quite sinusoidal. But in another sense the results are encouraging because we are comparing \( h_2(\xi) \) obtained by two entirely different methods. The ray theory results were obtained from the phase-front curvature in the near zone whereas I obtained the wave theory results from analyzing the far-zone rainbow. I obtained \( h_{ave} = 4.643 \) 972 for the elliptical-cross-section perturbation-theory data of Fig. 2, in comparison with \( h_2 = 4.643 \) 188 obtained using exact Rayleigh–Debye theory for a circular-cross-section cylinder having \( x_{ave} = 1000.05 \).

I verified that the difference between ray theory and first-order perturbation theory for \( h_2(\xi) - h_{ave} \) in Fig. 2 was not due to the perturbation theory calculation being beyond its upper limit for accuracy for an elliptical surface shape by decreasing \( \epsilon \) by an order of magnitude, repeating the calculation, and obtaining the same results. Similarly, I verified that the differences in Fig. 2 were not due to numerical inaccuracies in determining the intensity minima \( \theta_a^R \) and \( \theta_b^R \). An uncertainty \( \delta \theta_{ave} \) in degrees in either \( \theta_a^R \) or \( \theta_b^R \) produces an uncertainty

\[
\delta h_2 \sim (0.029 923) \delta \theta_{ave} (h_2 x_{ave})^{2/3},
\]

in \( h_2 \) and an uncertainty

\[
\delta \theta^R \sim \delta \theta_{ave}.
\]

in \( \theta^R \). When \( \delta \theta_{ave} = 0.000 01° \), we find that \( \delta h_2 \sim 0.000 086 \), which is approximately 2.2% of the amplitude of oscillation of \( h_2(\xi) \) in Fig. 2, and \( \delta \theta^R \sim 0.000 01° \), which is approximately 0.076% of the amplitude of oscillation of \( \theta^R(\xi) \) in Fig. 3. These uncertainties are negligible. Thus I conclude that the difference between the first-order perturbation theory with Airy theory modeling of the supernumerary intensity minima and the ray theory results for \( h_2(\xi) - h_{ave} \) are due to the fact that the wave theory intensity is not perfectly fit by the square of an Airy integral and, in particular, that the intensity minima in wave theory do not correspond exactly to any of the Airy integral of Eq. (7).

The sensitivity analysis described in the preceding paragraph also indicates that \( h_2 \) is approximately 29 times more sensitive (i.e., 2.2%/0.076% \( \sim 29 \)) to any inaccuracy in the determination \( \theta_a^R \) and \( \theta_b^R \) than is the position of the \( p = 2 \) rainbow. A consequence of this is that a sizable error in either the calculation or the measurement of \( h_2 \) can occur without greatly affecting the calculated or measured value of the rainbow angle. Said a different way, with Eqs. (58) written as

\[
\theta_2^R(\xi) = \theta_a^R(\xi) - 2.338 107 (180/\pi) h_2(\xi)^{1/3} / x_{ave}^{2/3},
\]

(63)
approximately 97% of the $\xi$ dependence of $\theta_2^R$ is produced by the $\xi$ dependence of $\theta_{\text{ave}}^{\text{min}}$ when $x_{\text{ave}} = 1000.0$ whereas only approximately 3% is produced by the $\xi$ dependence of $h_2$. This is why the comparison between first-order perturbation theory and ray theory for the rainbow angle was so good in Fig. 3. These differing sensitivities also have great significance for my experimental measurement of $h_2(\xi)$ and $h_3(\xi)$ to be reported separately.

We have already seen that use of the lowest-order CAM correction to Airy theory to model the intensity minima of the $p = 2$ rainbow of a circular-cross-section cylinder eliminated most of the difference in $\theta_2^R$ and $h_2$ between first-order perturbation theory and ray theory. We do not obtain the same degree of success, however, for an elliptical-cross-section cylinder. Use of CAM modeling of the rainbow intensity minima in Eqs. (58), the results for $\theta_2^R(\xi) - \theta_{\text{ave}}^{\text{min}}$ are identical to those of Airy rainbow modeling. Also, we obtain $\theta_{\text{ave}}^{\text{R}} = 137.914^\circ$ and $h_{\text{ave}} = 4.954$ 793 for the elliptical-cross-section cylinder which are extremely close to the CAM theory results $\theta_2^R = 137.914$ 91$^\circ$ and $h_2 = 4.954$ 462 for the circular-cross-section cylinder discussed above. But the amplitude of oscillation of $h_2(\xi) - h_{\text{ave}}$ by use of CAM modeling of the rainbow minima (with the coefficients $u_1$ and $v_0$ for a sphere) grows to 28% larger than that for ray theory. The CAM results are shown as the solid circles in Fig. 2. These results illustrate that, although CAM modeling of the rainbow minima improves the absolute value of $h_2$ for an elliptical-cross-section cylinder, it does not improve the $\xi$ dependence of $h_2$. If, however, the CAM formalism of Refs. 17–19 were to be extended to elliptical-cross-section cylinders so that the $\xi$ dependence of $u_1$ is included, the CAM predictions would again presumably be closer to ray theory than the predictions of first-order perturbation theory with Airy modeling of the rainbow minima.

5. Conclusion

The feature of the rainbow that has historically received the most attention by theorists is the far-zone rainbow scattering angle $\theta_p^R$. But both the ray theory and the first-order perturbation wave theory calculations described here indicate that the supernumerary spacing parameter $h_p$ is a much more sensitive and delicate feature of the rainbow than is $\theta_p^R$. Using ray theory I found that $h_p(\xi)$ oscillates sinusoidally as a function of $\xi$ for a small cylinder ellipticity $\epsilon$, just as the rainbow angle did in M"obius theory. Although I was not able to obtain an analytic expression for $h_p(\xi)$ for $\epsilon \ll 1$, a good approximation in this regime is given by approximations (16) and (17) for $p = 2$ and $1.25 \leq n \leq 1.7$. As $\epsilon$ increases, higher Fourier components become important, and the oscillation is no longer exactly sinusoidal.

In the first-order perturbation version of wave theory, we found in Eqs. (34) that each Fourier component of the cylinder’s surface shape for $q \geq 1$ couples two different incident partial waves to a single interior or scattered partial wave. The Debye-series decomposition of the perturbation-theory partial-wave scattering amplitudes of Eqs. (45) and (47) illustrates the physical mechanism of the partial-wave coupling in a clear way. The coupling occurs at any of the interactions of the multipole waves with the cylinder surface. The partial-wave-changing Fresnel coefficients for reflection and transmission are given in Eqs. (35) (38–40) (42), and (43). To determine $h_p(\xi)$ using first-order perturbation theory, we also had to model the minima of the supernumerary interference pattern using either Airy theory or the lowest-order CAM correction to Airy theory. We found that although CAM modeling for spheres produced a better agreement with ray theory for $h_2$ for a circular-cross-section cylinder and $h_{\text{ave}}$ for an elliptical-cross-section cylinder, the Airy modeling produced a better agreement for the oscillation of $h_2(\xi)$. We also found that reasonably large uncertainties can occur in $h_p(\xi)$ without seriously affecting the calculated value of the rainbow angle $\theta_p^R(\xi)$. Conversely, the angles of the supernumerary maxima or minima need to be measured with extreme accuracy to determine $h_p(\xi)$, which strongly affects experiments that attempt to measure $h_p(\xi)$.

This research was supported in part by the National Aeronautics and Space Administration under grant NCC-3-521.

References


36. Ref. 11, Table 10.13, p. 478.