A note on the green dyadic calculation of the decay rates for admolecules
at multiple planar interfaces

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The Green dyadic formulation for calculating classical decay rates of admolecules at multiple planar
interfaces first published by Chance, Prock and Silbey is reexamined. It is pointed out that, for the
case of fluorescing molecules sandwiched between a system of super- and substrate interfaces, the
original formalism requires significant modifications in order to lead to results consistent with those
obtained from the Sommerfeld theory. © 1999 American Institute of Physics.
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I. INTRODUCTION

The study of molecular fluorescence at solid interfaces has been active for the last 2 decades since the first experi-
mental works of Kuhn and Drexhage done in the early 1970’s. Theoretically, it has been found that one of the most
simple and direct approaches is to follow a phenomenologi-

cal model solving the electrodynamics of an emitting mo-

cular dipole in the vicinity of the interfaces. This model
can provide both frequency shifts and decay rates for the
admolecules can provide both frequency shifts and decay rates for the
lecular dipole in the vicinity of the interfaces. This model
provide both frequency shifts and decay rates for the
ADMolecules (normalized to the free molecular decay rate),
for example, yielding results in agreement with experimental
results as well as quantum mechanical calculations. Among
other contributors, Chance, Prock and Silbey (CPS) have
cleverly applied the radiation antenna theory of Sommerfeld to
this problem and showed that classical electrodynamics
alone can account for most of the experimental observations.
A momentous review article was compiled by CPS in the
late 1970’s summarizing the complete status of the subject at
that time. Over the last 20 years, this article has often been
quoted and used by people working in the field, experimental-
ists and theorists alike. Application of this theory has also
gone beyond fluorescence to other optical phenomena at in-
terfaces as in Ref. 6 and 7. It is also in this review article that
the Sommerfeld method was first generalized to the case of
multiple planar interfaces using the dyadic Green’s function
formulation. This generalization includes both cases (i) where
the layer of fluorescent dyes is deposited on the top of a
stratified multilayer system and (ii) where the layer is sand-
wiched between two such multilayer systems.

It is the purpose of this paper to point out that, in the
original CPS formulation for case (ii) above, the choice of
the dyadic eigenfunctions was not appropriately made and
significant modifications are necessary to obtain a consistent
Green dyadic theory for this case. The correct result will be
presented in two different but equivalent formulations and
will be shown to lead back to well-known results from the
Sommerfeld theory for the simple case with the dyes sand-
wiched between only one superstrate and one substrate
medium.

II. THE CPS FORMULATION

To be clear and self-contained, let us first recapitulate
the main results from the CPS article. For harmonic currents
and fields, the dyadic Green formulation of Ref. 4 has the
standard Green’s function solution (in SI units):

\[ E(R) = i \omega \mu \int G(R,R') \cdot J(R') dV(R') \]

(1)

where \( \mu \) is the magnetic permeability. For simplicity, we will
consider in this paper only the case with one superstrate and
one substrate confining the source in the gap as depicted in
Fig. 1. More details on the case with a large number of layers
and generalization to the case with gradient index media will
be presented in a forthcoming paper.

Let \( G_0 \) denote the Green dyadic for the source field and
\( G_i (i=1,2,3) \) denote those for the scattered fields in the three
media. Thus according to Ref. 4, one obtains

\[ G_0(R,R') = \frac{-1}{k_1^2} \frac{\delta(R-R')}{2\pi} + \frac{i}{4\pi} \int_0^{+\infty} d\lambda \sum_{n=0}^{+\infty} \frac{2 - \delta_n}{\lambda h_1(\lambda)} \sum_{j=0}^{1} \left[ M_{jn}(n+h_1)M_{jn}(n+h_1) + N_{jn}(n+h_1)N_{jn}(n+h_1) \right] \delta(z') \]

(2)

\[ G_1(R,R') = \frac{i}{4\pi} \int_0^{+\infty} d\lambda \sum_{n=0}^{+\infty} \frac{2 - \delta_n}{\lambda h_1(\lambda)} \sum_{j=0}^{1} \left[ f_j M_{jn}(n+h_1) + f_j' M_{jn}(n+h_1) \right] \delta(z') \]

(3)
where $\mathbf{M}$ and $\mathbf{N}$ are given in cylindrical coordinates by:

$$M_{j\alpha}(h) = e^{ihz} \left[ \frac{n J_n(\lambda r)}{r} \sin \left( \frac{j \pi}{2} - n \varphi \right) \hat{r} - \frac{\partial J_n(\lambda r)}{\partial r} \right] \times \cos \left( \frac{j \pi}{2} - n \varphi \right) \hat{\phi},$$

$$N_{j\alpha}(h) = \frac{e^{ihz}}{k} \left[ ih \frac{\partial J_n(\lambda r)}{\partial r} \cos \left( \frac{j \pi}{2} - n \varphi \right) \hat{r} + i n h J_n(\lambda r) \right] \cos \left( \frac{j \pi}{2} - n \varphi \right) \hat{\phi} + \lambda^2 J_n(\lambda r) \cos \left( \frac{j \pi}{2} - n \varphi \right) \hat{z}.$$

with $J_n$ the Bessel function of the first kind and $h_i(\lambda) = \sqrt{k_i^2 - \lambda^2}$, where the square root is taken to have positive real part.

According to Ref. 4, requirement of continuity of transverse field components at interfaces $z=0$ and $z=z_0$ leads to the following systems where $e_j = e^{ih_j}z_0$:

$$\begin{align*}
\mathbf{G}_2(\mathbf{R}, \mathbf{R}') &= \frac{i}{4\pi} \int_0^{+\infty} d\lambda \sum_{n=0}^{+\infty} \frac{2 - \delta_n}{\lambda h_1(\lambda)} \\
&\times \sum_{j=0}^{1} \left[ c_2 M_{j\alpha}(h_2) M'_{j\alpha}(h_1) + f_2 N_{j\alpha}(h_2) N'_{j\alpha}(h_1) \right], \\
\mathbf{G}_3(\mathbf{R}, \mathbf{R}') &= \frac{i}{4\pi} \int_0^{+\infty} d\lambda \sum_{n=0}^{+\infty} \frac{2 - \delta_n}{\lambda h_1(\lambda)} \\
&\times \sum_{j=0}^{1} \left[ c_3 M_{j\alpha}(h_3) M'_{j\alpha}(h_1) + f_3 N_{j\alpha}(h_3) N'_{j\alpha}(h_1) \right],
\end{align*}$$

where $\mathbf{M}$ and $\mathbf{N}$ are given in cylindrical coordinates by:

$$\begin{align*}
\mathbf{M}_{j\alpha}(h) &= e^{ihz} \left[ \frac{n J_n(\lambda r)}{r} \sin \left( \frac{j \pi}{2} - n \varphi \right) \hat{r} - \frac{\partial J_n(\lambda r)}{\partial r} \right] \\
&\times \cos \left( \frac{j \pi}{2} - n \varphi \right) \hat{\phi},
\end{align*}$$

$$\begin{align*}
\mathbf{N}_{j\alpha}(h) &= \frac{e^{ihz}}{k} \left[ ih \frac{\partial J_n(\lambda r)}{\partial r} \cos \left( \frac{j \pi}{2} - n \varphi \right) \hat{r} + i n h J_n(\lambda r) \right] \\
&\times \sin \left( \frac{j \pi}{2} - n \varphi \right) \hat{\phi} + \lambda^2 J_n(\lambda r) \cos \left( \frac{j \pi}{2} - n \varphi \right) \hat{z}.
\end{align*}$$

or in matrix form $\mathbf{A}e = r_e$, and

$$\begin{pmatrix}
-h_1/k_1 & h_1/k_1 & 0 & h_3/k_3 \\
k_1 & k_1 & 0 & -k_3 \\
-h_1/k_1 e_1 & h_1 e_1/k_1 & -h_2 e_2/k_2 & 0 \\
k_1 e_1 & k_1 e_1 & -k_2 e_2 & 0
\end{pmatrix} = \begin{pmatrix}
h_1/k_1 \\
-k_1 \\
-e_1 h_1/k_1 \\
-e_1 k_1
\end{pmatrix},$$

or $\mathbf{B}e = r_f$. Solving Eqs. (8) and (9) yields the following:

$$\begin{pmatrix}
1 & 1 & 0 & -1 \\
-h_1 & h_1 & 0 & h_3 \\
1 & e_1 & -e_2 & 0 \\
h_1 e_1 & h_1 e_1 & -h_2 e_2 & 0
\end{pmatrix} \begin{pmatrix}
c_1 \\
c'_1 \\
c_2 \\
c'_2
\end{pmatrix} = \begin{pmatrix}
-1 \\
-h_1 \\
-e_1 \\
-h_1 e_1
\end{pmatrix},$$

where

$$R_{ij}^0 = \frac{e_i h_j - e_j h_i}{e_i h_i + e_j h_j}$$

and

$$R_{ij}^1 = \frac{h_i - h_j}{h_i + h_j}.$$
A. Solution by expanding the solution space

In reviewing the problem, we found that Eqs. (8) and (9) do not satisfy the boundary conditions at the interfaces and that no solution could be found once the constraints of forms (3) (4), and (5) were imposed. A necessary remedy is enlargement of the solution space to the point where the boundary conditions can be satisfied. For instance, the expression for \( G_1 \) contains dyadic products \( M_{jnk}(-h_1)M'_{jnk}(-h_1) \) and \( M_{jnk}(+h_1)M'_{jnk}(+h_1) \) but not \( M_{jnk}(-h_1)M'_{jnk}(+h_1) \) and \( M_{jnk}(+h_1)M'_{jnk}(-h_1) \) which are equally valid. It turns out that the correct solution from this approach has already been worked out in the electrical engineering literature.10 The general solutions for the scattering fields are given by

\[
G_1(\mathbf{R}, \mathbf{R}') = \frac{i}{4\pi} \int_0^{+\infty} d\lambda \sum_{n=0}^{+\infty} \frac{2-\delta_n}{\lambda h_1(\lambda)} \left\{ [c_1 M_{jnk}(-h_1) + c'[M_{jnk}(h_1)]M'_{jnk}(h_1) \right. \\
+ [a_1 M_{jnk}(-h_1) + a'[M_{jnk}(h_1)]M'_{jnk}(h_1) + [f_1 N_{jnk}(-h_1) + f'[N_{jnk}(h_1)]N'_{jnk}(h_1) \\
+ [b_1 N_{jnk}(-h_1) + b'[N_{jnk}(h_1)]N'_{jnk}(-h_1)] \right\}. 
\]

\[
G_2(\mathbf{R}, \mathbf{R}') = \frac{i}{4\pi} \int_0^{+\infty} d\lambda \sum_{n=0}^{+\infty} \frac{2-\delta_n}{\lambda h_1(\lambda)} \left\{ [c_2 M_{jnk}(h_2)M'_{jnk}(h_1) + f_2 N_{jnk}(h_2)N'_{jnk}(h_1) + a_2 M_{jnk}(h_2)M'_{jnk}(h_1) \right. \\
+ b_2 N_{jnk}(h_2)N'_{jnk}(-h_1)] \right\}. 
\]

\[
G_3(\mathbf{R}, \mathbf{R}') = \frac{i}{4\pi} \int_0^{+\infty} d\lambda \sum_{n=0}^{+\infty} \frac{2-\delta_n}{\lambda h_1(\lambda)} \left\{ [c_3 M_{jnk}(-h_3)M'_{jnk}(h_1) + f_3 N_{jnk}(-h_3)N'_{jnk}(h_1) \right. \\
+ b_3 N_{jnk}(-h_3)N'_{jnk}(-h_1)] \right\}. 
\]

By imposing the appropriate boundary conditions, the expansion coefficients can finally be obtained as

\[
c_s = \begin{bmatrix} c_1 \\ c_1' \\ c_2 \\ c_2' \\ c_3 \\ c_3' \end{bmatrix}, \quad a_s = \begin{bmatrix} a_1 \\ a_1' \\ a_2 \\ a_2' \\ a_3 \\ a_3' \end{bmatrix}, \quad f_s = \begin{bmatrix} f_1 \\ f_1' \\ f_2 \\ f_2' \end{bmatrix}, \quad b_s = \begin{bmatrix} b_1 \\ b_1' \\ b_2 \\ b_2' \end{bmatrix}. 
\]

Using the above solution, we can write out the Green’s functions as

\[
G_1(\mathbf{R}, \mathbf{R}') = \frac{i}{4\pi} \int_0^{+\infty} d\lambda \sum_{n=0}^{+\infty} \frac{2-\delta_n}{\lambda h_1(\lambda)} \left\{ \frac{1}{1-e^{2}\mathbf{R}_{12}'\mathbf{R}_{13}'} \right. \\
\times \left[ e^{2}\mathbf{R}_{12}'\mathbf{M}'(\mathbf{R}_{13}'\mathbf{M}'+\mathbf{M}') + \mathbf{R}_{13}'\mathbf{M}'(\mathbf{N}'+\mathbf{M}'+\mathbf{M}') \right. \\
+ \left. e^{2}\mathbf{R}_{12}'\mathbf{M}'(\mathbf{N}'-\mathbf{N}') \right] \right\}, 
\]
where \( \mathbf{M}^\pm = \mathbf{M}'(\pm h_1) \) and \( \mathbf{N}^\pm = \mathbf{N}'(\pm h_1) \). We have checked that the solution given by Eqs. (15)–(17) does satisfy the numerical test described above.

### B. Solution by reassociation

An alternative approach, which might be called “‘reassociation’”, is to introduce explicitly the source \( \mathbf{J} \) into the dyadic expansion. We replaced typical products such as \( (\mathbf{M}^\prime \mathbf{M}')\mathbf{J} \) with the equivalent product \( \mathbf{M}(\mathbf{M}^\prime \cdot \mathbf{J}) \) reducing the product on the right to a complex scalar. An additional small step then leads to the realization that \( c(\mathbf{M}^\prime \mathbf{M}')\mathbf{J} \) can be replaced by \( c\mathbf{M} \), where \( \mathbf{M}^\prime \cdot \mathbf{J} \) scalar has been absorbed into the \( c \). We will see at the end that \( \mathbf{J} \) can be factored from both sides of the resulting equations leading to expressions for \( \mathbf{G} \), independent of the source, as they must be. Following the logic given above, we have:

\[
\mathbf{G}_1(\mathbf{R}, \mathbf{R}^\prime) = \frac{i}{4\pi} \int_0^{+\infty} \frac{d\lambda}{\lambda} \sum_{n=0}^{+\infty} \sum_{j=0}^{+\infty} \frac{2-\delta_n}{\lambda h_1(\lambda)} c_1 \mathbf{M}_{jn\lambda}(-h_1) + c'_1 \mathbf{M}_{jn\lambda}(h_1) + f_1 \mathbf{N}_{jn\lambda}(-h_1) + f'_1 \mathbf{N}_{jn\lambda}(h_1),
\]

\[
\mathbf{G}_2(\mathbf{R}, \mathbf{R}^\prime) = \frac{i}{4\pi} \int_0^{+\infty} \frac{d\lambda}{\lambda} \sum_{n=0}^{+\infty} \sum_{j=0}^{+\infty} \frac{2-\delta_n}{\lambda h_1(\lambda)} c_2 \mathbf{M}_{jn\lambda}(h_2) + f_2 \mathbf{N}_{jn\lambda}(h_2),
\]

\[
\mathbf{G}_3(\mathbf{R}, \mathbf{R}^\prime) = \frac{i}{4\pi} \int_0^{+\infty} \frac{d\lambda}{\lambda} \sum_{n=0}^{+\infty} \sum_{j=0}^{+\infty} \frac{2-\delta_n}{\lambda h_1(\lambda)} c_3 \mathbf{M}_{jn\lambda}(-h_3) + f_3 \mathbf{N}_{jn\lambda}(-h_3),
\]

where the \( c \) and \( f \) coefficients are functions of current density \( \mathbf{J} \) as well as position, \( j, n \) and \( \lambda \). This approach allows us to work directly with electric and magnetic field values in applying the boundary conditions at the interfaces. The resulting eight equations decouple into two matrix systems:

\[
\mathbf{Ac} = \left[ -\mathbf{M}'(h_1) \quad h_1 \mathbf{M}'(h_1) \quad -e_2 \mathbf{M}'(-h_1) \quad -h_1 e_1 \mathbf{M}'(-h_1) \right] \mathbf{J}',
\]

\[
\mathbf{Bf} = \left[ \frac{h_1}{k_1} \mathbf{N}'(h_1) \quad -k_1 \mathbf{N}'(h_1) \quad -\frac{h_1 e_1}{k_1} \mathbf{N}'(-h_1) \quad -k_1 e_1 \mathbf{N}'(-h_1) \right] \cdot \mathbf{J}'.
\]

where \( \mathbf{A} \) and \( \mathbf{B} \) are the same matrices defined above in Eqs. (8) and (9). The matrix solution of Eqs. (21) and (22) leads to:

\[
\begin{bmatrix}
  c_1 \\
  c'_1 \\
  c_2 \\
  c_3
\end{bmatrix} = \frac{1}{1-e^{-i \frac{2\pi}{\lambda R_{12}} R_{13} \mathbf{M}' + \mathbf{M}^\prime}} \cdot \mathbf{J},
\]

\[
\begin{bmatrix}
  f_1 \\
  f'_1 \\
  f_2 \\
  f_3
\end{bmatrix} = \frac{1}{1-e^{-i \frac{2\pi}{\lambda R_{12}} R_{13} \mathbf{N}' - \mathbf{N}^\prime}} \cdot \mathbf{J}.
\]

We can now insert Eqs. (23) and (24) into Eqs. (18)–(20). Since an arbitrary \( \mathbf{J} \) then appears on both sides of the result, we can factor out \( \mathbf{J} \), yielding results in complete agreement with Eqs. (15)–(17). More details on this approach and the equivalence between the two methods will be provided in a forthcoming paper.8

### IV. CALCULATION OF DECAY RATES

According to the classical phenomenological approach of CPS, the normalized decay rate of the admolecule can be obtained in terms of the imaginary part (J) of the reflected field at the dipole site as:

\[
\hat{b} = \frac{b}{b_0} = 1 + \frac{6 \pi \epsilon_0 q n^2}{p_0 k^3} \gamma(E_0),
\]

(25)
where \( q \) is the intrinsic quantum yield and \( k_1 = n_1 \omega / c \), with \( n_1 \) the real refractive index of the medium containing the dipole. We show below that the dyadics given in Eqs. (15)–(17) can indeed lead back to the correct results for \( \hat{b} \) for a molecule confined as in the geometry of Fig. 1.

We first consider the problem of a vertically oriented dipole with moment \( p_{0} \hat{z} e^{-i \omega t} \) at the source position \( d \hat{z} \) between two interfaces at \( z = 0 \) and \( z = s + d = z_0 \). The current will be given by

\[
\mathbf{J} = -i \omega p_{0} \hat{z} e^{-i \omega t} \delta(\mathbf{R}' - d \hat{z}).
\]

Inserting this \( \mathbf{J} \) into Eqs. (1) and using Eq. (15) yields

\[
\begin{align*}
\mathbf{E}_z'(d \hat{z}) &= \left\{ \omega^2 \mu_0 \mu_1 p_{0} e^{-i \omega t} \right\} \hat{z} \cdot \mathbf{G}_1(d \hat{z}, d \hat{z}) \hat{z} \\
&= \frac{1}{2 \pi} \int_{- \infty}^{+ \infty} \sum_{\lambda = - \infty}^{+ \infty} \lambda h_1(\lambda) \\
&\times \sum_{j = 0, 1} \left\{ \frac{1}{1 - e^{2 R_{12} R_{13}}} [e_1^2 R_{12} N^-(R_{13} N^- - N^+) + R_{13} N^+ (e_1^2 R_{12} N^- - N^+)] \right\} \hat{z} d \lambda \\
&= \frac{1}{2 \pi} \int_{- \infty}^{+ \infty} \lambda h_1(\lambda) \\
&\times \sum_{j = 0, 1} \left\{ \frac{1}{1 - e^{2 R_{12} R_{13}}} [e_1^2 R_{12} N^-(R_{13} N^- - N^+) + R_{13} N^+ (e_1^2 R_{12} N^- - N^+)] \right\} \hat{z} d \lambda \\
&= \frac{1}{2 \pi} \int_{- \infty}^{+ \infty} \lambda h_1(\lambda) \\
&\times \sum_{j = 0, 1} \left\{ \frac{1}{1 - e^{2 R_{12} R_{13}}} [e_1^2 R_{12} N^-(R_{13} N^- - N^+) + R_{13} N^+ (e_1^2 R_{12} N^- - N^+)] \right\} \hat{z} d \lambda.
\end{align*}
\]

where \( e_d = e^{i \phi_d(\lambda)} \) and \( e_x = e^{i \phi_x(\lambda)} \) with \( s + d = z_0 \). We have also employed the identity:

\[
\frac{2 \chi_y - x - y}{1 - x y} = \frac{1 - \chi(1 - y)}{1 - x y} - 1.
\]

Inserting Eq. (27) into Eq. (25), we obtain

\[
\begin{align*}
\hat{b}_z &= 1 + \frac{6 \pi \epsilon_0 \omega q n_1^2}{p_0 k_1^2} \left[ \frac{3}{2 \pi} \int_{- \infty}^{+ \infty} \lambda h_1(\lambda) \\
&\times \left\{ \frac{1 - R_{12} R_{13}}{1 - e^{2 R_{12} R_{13}}} \right\} \hat{z} d \lambda \right] \\
&= 1 - q + \frac{3 q}{2 \pi} \\
&\times \int_{- \infty}^{+ \infty} \lambda h_1(\lambda) \\
&\times \left\{ \frac{1 - R_{12} R_{13}}{1 - e^{2 R_{12} R_{13}}} \right\} \hat{z} d \lambda,
\end{align*}
\]

where we have used

\[
\int_{\lambda = 0}^{\lambda = \sqrt{\chi_1^2 - \chi_2^2}} \frac{\lambda^3 d \lambda}{\sqrt{\chi_1^2 - \chi_2^2}} = \frac{2}{3} \lambda^3.
\]

The result in Eq. (28) is equivalent to Eq. (2.47) of Ref. 4 using the transformation: \( u = \lambda / k_1 \). In the case of a trivial interface between regions 1 and 2, that is, \( \epsilon_1 = \epsilon_2 \) implying \( R_{12} = 0 \), from Eq. (28) we have

\[
\hat{b}_z = 1 - \frac{3 q}{2 \pi} \int_{- \infty}^{+ \infty} \lambda h_1(\lambda) \\
&\times \left\{ \frac{1 - R_{12} R_{13}}{1 - e^{2 R_{12} R_{13}}} \right\} \hat{z} d \lambda.
\]

which is identical to Eq. (2.17) of Ref. 4.

We next consider the problem of a horizontally oriented dipole with moment \( p_{0} \hat{x} e^{-i \omega t} \) at the source position \( d \hat{z} \) between the same interfaces. The current density is then

\[
\mathbf{J} = -i \omega p_{0} \hat{x} e^{-i \omega t} \delta(\mathbf{R}' - d \hat{z}).
\]

Again, inserting \( \mathbf{J} \) into Eq. (1) and integrating yields

\[
\mathbf{E}_z'(d \hat{z}) = \left\{ \omega^2 \mu_0 \mu_1 p_{0} e^{-i \omega t} \right\} \hat{z} \cdot \mathbf{G}_1(d \hat{z}, d \hat{z}) \hat{x}.
\]

To proceed further, we note that

\[
\hat{x} \cdot \mathbf{M}^z(d \hat{z}) = \begin{cases} \lambda \frac{e^{\pm i \phi} d \lambda}{2} & \text{if } n = 1 \text{ and } j = 1, \\ 0 & \text{otherwise} \end{cases}
\]

and

\[
\hat{x} \cdot \mathbf{N}^z(d \hat{z}) = \begin{cases} \pm i \lambda h_1 \frac{e^{ \pm i \phi} d \lambda}{2 k_1} & \text{if } n = 1 \text{ and } j = 0, \\ 0 & \text{otherwise} \end{cases}
\]

Insertion of the above dot products into Eq. (15) yields
\[ \hat{x} \cdot \mathbf{G}_1(d\hat{x}, d\hat{z}) = \frac{i}{8\pi} \int_{\lambda=0}^{+\infty} \frac{1}{1-e^{-iR_{12}R_{13}}(2e_i^2R_{12}^4R_{13}^4 + e_i^2R_{12}^4 + e_i^2R_{13}^4)}
\[ + \frac{h^2_{\perp}(\lambda)}{k_1^2(1-e^{-iR_{12}R_{13}})(2e_i^2R_{12}^4R_{13}^4 + e_i^2R_{12}^4 + e_i^2R_{13}^4)} \frac{\lambda}{h_1(\lambda)} ] d\lambda, \]

where we have again used the identity:

\[ \frac{2xy+x+y}{1-xy} = \frac{(x+1)(y+1)}{1-xy} - 1. \]

Using \( k_1^2 = \omega^2 \varepsilon_0 \varepsilon_1 \mu_0 \mu_1 \), we finally have

\[ E_1(d\hat{z}) = \frac{p_0 e^{-i\omega t}}{8\pi \varepsilon_0 \varepsilon_1} \int_{\lambda=0}^{+\infty} \left[ k_1^2 \left( \frac{(e_i^2R_{12}^4R_{13}^4 + e_i^2R_{12}^4 + e_i^2R_{13}^4)}{1-e^{-iR_{12}R_{13}}(2e_i^2R_{12}^4R_{13}^4 + e_i^2R_{12}^4 + e_i^2R_{13}^4)} - 1 \right) + h^2_{\perp}(\lambda) \left( \frac{(e_i^2R_{12}^4R_{13}^4 + e_i^2R_{12}^4 + e_i^2R_{13}^4)}{1-e^{-iR_{12}R_{13}}(2e_i^2R_{12}^4R_{13}^4 + e_i^2R_{12}^4 + e_i^2R_{13}^4)} - 1 \right) \right] \frac{\lambda}{h_1(\lambda)} ] d\lambda. \]

Inserting Eq. (32) into Eq. (25), we obtain

\[ \hat{h}_1 = 1 + \frac{6 \pi \varepsilon_0 q n_1^2}{p_0 k_1^2} \frac{p_0}{8\pi \varepsilon_0 \varepsilon_1} \int_{\lambda=0}^{+\infty} \left[ k_1^2 \left( \frac{(e_i^2R_{12}^4R_{13}^4 + e_i^2R_{12}^4 + e_i^2R_{13}^4)}{1-e^{-iR_{12}R_{13}}(2e_i^2R_{12}^4R_{13}^4 + e_i^2R_{12}^4 + e_i^2R_{13}^4)} - 1 \right) + h^2_{\perp}(\lambda) \left( \frac{(e_i^2R_{12}^4R_{13}^4 + e_i^2R_{12}^4 + e_i^2R_{13}^4)}{1-e^{-iR_{12}R_{13}}(2e_i^2R_{12}^4R_{13}^4 + e_i^2R_{12}^4 + e_i^2R_{13}^4)} - 1 \right) \right] \frac{\lambda}{h_1(\lambda)} ] d\lambda \]

where we have used

\[ \int_{\lambda=0}^{+\infty} \frac{(2k^2-\lambda^2)\lambda d\lambda}{\sqrt{k^2-\lambda^2}} = \frac{4}{3} k^3. \]

The result in Eq. (33) is equivalent to Eq. (2.48) of Ref. 4 using the transformation: \( u = \lambda/k_1 \). In the case of a trivial interface between regions 1 and 2, that is \( \varepsilon_1 = \varepsilon_2 \), implying that \( R_{12} = R_{13} = 0 \), we have from Eq. (33)

\[ \hat{h}_1 = 1 + \frac{3q}{4k_1^2} \int_{\lambda=0}^{+\infty} e_i^2 \left[ k_1^2 h^2_{\perp}(\lambda) R_{13}^4 \lambda d\lambda / h_1(\lambda) \right]. \]

which is identical to Eq. (2.29) of Ref. 4. Thus our Green dyadics in Eqs. (15)–(17) indeed reproduce the correct results for the decay rates obtained by the generalization of the Sommerfeld method.4

V. CONCLUSION

The dyadic Green’s function solution to the double mirror problem in Ref. 4 was found to have a theoretical error. We have shown that the error can be corrected by extending and symmetrizing the solution form in order to satisfy boundary conditions as done in Ref. 10. Alternatively, reassociation of the dyadic product with current density leads to the same solution as can be seen by virtue of an isomorphism between formulations.8 Finally, the corrected dyadic Green’s function formulation can be used directly to calculate and verify decay rates calculated in Ref. 4 from the Sommerfeld theory for the case of an oscillating dipole positioned between interfaces. With either of the two approaches, generalization is straightforward to the case with an arbitrary number of multiple interfaces for both the substrate and superstrate.8,10 In addition, the green dyadlic formalism will also allow one to calculate an arbitrary source within the gap beyond that of an electric point dipole. The present formulation should be useful in these aspects.

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9. Note that the \( \delta \) function term is necessary for \( G_0 \) as explained in C. T. Tai, Proc. IEEE 61, 480 (1973).