for almost Lorentzian mixtures[35] has been suggested as a method of overcoming this difficulty. A far simpler approach is used here: at low temperatures, where the convergence is slow, the Lorentzian expressions,[32] modified[31,36] to yield the correct limiting behavior at full ionization, are used. They read

$$\sigma \equiv \sigma^{\perp} + i\sigma^{H} = \frac{4 e^2 n_e}{3 \sqrt{\pi} m_e} \int_o^{\infty} \frac{e^{-x} x^{3/2} dx}{\nu_e (kTx) + i \omega_e}$$ (23)

$$D_e^T \equiv D_e^{T\perp} + iD_e^{TH} = \frac{4 n_e kT}{3 \sqrt{\pi}} \int_o^{\infty} \frac{e^{-x}(\frac{5}{2} - x) x^{3/2} dx}{\nu_e (kTx) + i \omega_e}$$ (24)

$$\lambda_e' \equiv \lambda_e'^{\perp} + i\lambda_e'^{H} = \frac{4 n_e k^2 T}{3 \sqrt{\pi} m_e} \int_o^{\infty} \frac{e^{-x}(\frac{5}{2} - x)^2 x^{3/2} dx}{\nu_e (kTx) + i \omega_e}$$ (25)

with

$$\nu_e = \nu_{ea} + \nu_{ei}$$ (26)

$$\nu_{ea} = n_a g Q^{(1)}(E)$$ (27)

$$\nu_{ei} = A_{\sigma, D, \lambda} \frac{16}{3g^2} n_i b_o^2 (\frac{2\pi kT}{m_e})^{\frac{1}{2}} \ln \Lambda$$ (28)

$A_{\sigma} = 1.27$, $A_D = 2.70$ and $A_{\lambda} = 1.80$ are constants introduced[31,36] to insure that the above expressions become those of Spitzer and Härm at full ionization. λ_e' of Eq. (25) is not the true thermal conductivity of Eq. (18). Each component of λ_e is found from the relation

$$\lambda_e = \lambda_e' - \frac{e^2 (D_e^T)^2}{\sigma m_e T}$$ (29)

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Since the above relations approach the cut-off expressions near full ionization, it was decided to use them only at very low ionization and low magnetic fields where problems of convergence in the Sonine polynomial approximations are present. This step avoids the use of auxiliary functions necessary to ensure the accuracy of Eqs. (22)-(25) at all values of magnetic field,[31] and allows for the use of the more accurate shielded cross sections.

The third approximation was used otherwise for the electron properties. The expressions necessary have been given earlier in some detail,[29,30] and will not be repeated here.* The interpolation between the results from the two different methods was found to be straightforward and occurred, for example, around 6000°K for p = 1 atm. All ion species were considered in computing the electron properties, since it is a simple matter to include extra ions in the Chapman-Enskog expressions.

Inasmuch as the transport coefficients, particularly λ_r, depend critically on the thermodynamic properties of the mixture, some information is included here on the methods used for this aspect of the computations. An equation for n_e in terms of the various partition functions was obtained from the law of mass action for each ionizing reaction as in Ref. 37. The partition functions were evaluated by two techniques -- for the ions the first few low energy levels and some not in the principal series were taken from Ref. 38. The contribution from the remaining levels were estimated with the aid of the simple formula[27]

$$ Z_j' = \frac{2}{3} (2 J_c + 1) \left(\frac{z_C E_H}{\Delta E_\infty^j} \right)^{\frac{3}{2}} \exp \left(\frac{\Delta E_\infty^j - E_\infty^j}{k T} \right) \tag{30} $$

where J_c is the total angular momentum quantum number of the core and z_C is its charge, E_H is the hydrogen ionization potential, E_∞^j is the ionization potential of species j and ΔE_∞^j is the lowering of E_∞^j due to the microfields. The electronic excitation energy corresponding to the partition function of Eq. (30) can be shown to be

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* Complete formulas are given in Appendix B.

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