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Discharge-Pumped Phasers

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The possibility of realizing an efficient gaseous laser-beam-generating medium that utilizes coherently phased (*i.e.* "dressed") atoms for the active laser species, but that does not inherently require the use of external laser beams for pumping, is explored. From dressed-atom energy level diagrams, it is immediately apparent that continuous-wave (CW) optical gain at the two resonance frequencies ω_o and ω_o' of Λ-type coherently phased atoms can only result when such atoms are additionally excited via nonlinear photonic processes. Non-parametric excitation processes of this type are considered - specifically, multi-photon stimulated hyper-Raman scattering (SHRS) processes driven by fluorescence radiation generated in a continuous electrical discharge present within the vapor cell. It is shown that no CW gain could result from dressed-atom excitation occurring via three-photon SHRS under any circumstances. Four-photon SHRS is a conceptually possible pumping means for such a device, and rough calculations indicate that the amount of optical gain produced by this high-order process should be marginally sufficient to allow multiwatt CW laser operation to occur on atomic transitions with levels relatively positioned to form a "double-Λ" structure. However, to initiate operation of such a device would require injection into the laser optical cavity of an intense "starter" laser pulse. An experimental configuration is described for determining feasibility of the proposed laser device. In the suggested configuration, Cs-atom $6S_{1/2}-6P_{1/2}$ transitions form the double-Λ structure.

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I. INTRODUCTION

Over the past fifteen years or so, many theoretical and experimental studies have shown that radically new pathways for efficiently generating laser radiation can be accessed if use is made of coherently phased (*i.e.* "dressed") atoms and molecules for the active laser species [1]. An illustration of this statement is provided by a recently published experimental study [2]. In Ref. [2], co-propagation in ^{208}Pb vapor of three laser pulses (with Rabi frequencies Ω_c, Ω_p, and Ω_e), each resonant with a separate allowed transition in a double-Λ configuration, resulted in efficient parametric generation of a co-propagating fourth laser pulse Ω_h, with the frequency ω_h of the latter being that of the remaining transition in the double-Λ level structure. This generation scheme embodies several recently discovered physical principles that are not active in conventional four-wave mixing experiments. To begin with, the synchronously applied "strong" laser pulses Ω_c and Ω_p, acting on a pair of transitions that share a common upper level, coherently phase the atoms of the vapor, driving each atom into a superposition dark state which becomes non-absorbing at the frequencies ω_c and ω_p of the two applied pulses, thereby allowing the latter to propagate orders of magnitude further into the vapor than if the atoms were not coherently phased. This establishment of electromagnetically induced transparency (EIT) also creates an atomic coherence ρ_{12} having an absolute value close to ½, the maximum value theoretically possible. This large atomic coherence, in turn, acts as a local oscillator that can beat with other co-propagating laser pulses, such as the applied pulse at ω_e, producing sum and difference frequencies. In this case, it is light at the sum frequency $\omega_h = (\omega_p - \omega_c) + \omega_e$ that is exactly resonant and that becomes efficiently generated. Were the atoms not coherently phased, the generated light at ω_h would be very strongly attenuated. Here, however, with the effects of quantum interference and atomic coherence playing important roles, a remarkable redistribution of intensities and phases of the pulses at ω_e and ω_h automatically occurs, resulting in the additional establishment of EIT on both transitions. All four laser pulses then propagate stably together as "matched pulses" in a completely loss-free manner, with the asymptotically attained Rabi frequencies Ω_e and Ω_h satisfying the condition $\Omega_e / \Omega_h = \Omega_c / \Omega_p$. Having all four laser beams exactly resonant, of course, allows access to nonlinear mixing susceptibilities that are orders of magnitude greater than are available in conventional four-wave mixing experiments. In Ref. [2], this enabled significant $\Omega_e \to \Omega_h$ photon conversion to occur over very short interaction lengths, and contributed importantly to the high efficiency observed for this process. Additionally, when this new approach to laser beam generation is followed, it no longer becomes necessary to seek ways to phase match the beams, since the latter occurs automatically when all the beams are resonant.

In the scheme employed in Ref. [2], and (to our knowledge) in all other published schemes that have focused upon the use of coherently phased atoms or molecules for laser beam generation, one or more externally applied laser beams are required both to coherently phase the active species, and to supply photons that can be converted to those of an output beam generated at another frequency. This statement, of course, does not hold for the vast majority of standard laser types, *i.e.* those requiring that population inversions be present on lasing transitions (but not requiring that the active species be coherently phased). One can therefore reasonably inquire whether it might also be possible to realize a laser that is based upon coherently phased atoms or molecules, but that is pumped in some manner by incoherent light, so that application of external laser beams is not intrinsically required to power the system. In the present paper a specific device of this type is proposed, and its potential performance characteristics are analyzed.

One begins a description of the system here envisioned by specifying the laser medium to be a low-pressure (: 0.1-1 Torr) atomic vapor coherently phased by the passage of two co-propagating, continuous-wave (CW) laser beams Ω and Ω' through the medium. Each laser beam has a frequency that is exactly resonant with one member of a pair of atomic resonance-line transitions that share a common upper level - *i.e.* the participating transitions form a Λ-type structure. For reasons that will become apparent, there should also be another nearby upper level that is radiatively coupled to both lower Λ levels, *i.e.* the four levels should form a double-Λ structure. Although several atomic candidates can be considered for the active species in such a system, the use of ^{133}Cs atoms appears to be especially advantageous and, for the sake of definiteness, will be assumed in most of the calculations that follow.

It will be shown that a practicable CW laser pumped by incoherent light would be impossible to realize if the active laser species were coherently phased cascade- or V-type three-level atoms (Sec. IIA). However, it will also be shown in the present paper that a gas of coherently phased atoms possessing a double-Λ energy level structure can serve as the basis for such a device, if the atoms are additionally excited by a specific type of nonlinear photonic process, with the latter being driven entirely by intense incoherent light that is spectrally concentrated about the double-Λ resonance frequencies.

It is assumed that a continuously occurring low-pressure electrical discharge is maintained within the vessel that contains the atomic vapor, and that intense atomic resonance fluorescence results from electron impact collisions with the atoms. Since, in the model, it is this fluorescence which drives the gain-producing nonlinear photonic processes, it is of course necessary to assume that high fluorescence photon intensities are present in the region of the vapor cell traversed by the co-propagating laser beams Ω and Ω'. Some discussion of this assumption is required, however, since standard theories of continuously operating fluorescent lamp discharges indicate that, due to radiation trapping, most of the energy produced via electron impact with atoms of the gas resides in the medium as atomic excitation, with purely photonic (*i.e.* electromagnetic) energy representing only a tiny fraction of the stored energy density (Sec. IIIB).

The present paper is organized as follows. Throughout Sec. II, dressed-atom level diagrams for three-level atomic systems [3] are utilized to analyze in a relatively simple way various photonic processes by which dressed atoms can be optically excited, the aim here being to see if optical gain can thereby be produced at the bare-atom transition frequencies. In Sec. IIA, it is shown that any CW laser, in which optical gain results from a coherently phased three-level atomic gas being pumped by incoherent light, could only work with an active species possessing a Λ structure – *i.e.* that use of cascade- or V-type dressed-atom systems in this application can be ruled out.

A general argument is then presented in Sec. IIB to show why direct optical pumping via the *linear* absorption spectrum of a coherently phased Λ gas cannot by itself produce optical gain on the Λ transitions. One thus is led to explore whether amplification of the propagating CW laser beams might result if the dressed atoms are excited nonlinearly.

Nonlinear excitation of dressed atoms via *non-parametric* (*i.e.* Raman-type) processes is explored in Sec. IIC. Here, one is essentially restricted to consider the class of multi-photon processes known as stimulated hyper-Raman scattering (SHRS). It becomes immediately apparent that dressed-atom excitation via the canonical *three*-photon form of SHRS could not result in the production of optical gain under any circumstances. *Four*-photon SHRS is considered next. It is shown that optical gain conceptually could be realized in the proposed laser with this process employed as a pumping mechanism. Moreover, estimates made in Sec. IID of the optical gain that could result from this high order process suggest that such gain would be at least marginally sufficient to sustain powerful CW laser operation in the case of a specific double-Λ pumping scheme involving Cs atoms. In this scheme, the effective pump power for lasing at one Cs double-Λ transition pair is the discharge-produced fluorescence occurring at the other. In principle, CW lasing could also occur in this manner simultaneously on all four double-Λ transitions. This potentially very efficient regime would be entirely driven by fluorescence pump energy occurring at all four double-Λ transitions. In Sec. IID, the highly relevant question of whether or not multi-photon SHRS processes can be driven by purely stochastic light is also briefly considered.

For the sake of completeness, *parametric*-type nonlinear photonic processes are also briefly considered in Sec. IIE as potential pumping mechanisms for the proposed device. However, on general grounds, it is shown that such processes would be ineffective in the present case.

In Sec. III experimental aspects of the proposed laser device are discussed. A generic type of vessel that enables a long column of hot reactive gas such as Cs vapor to be stably contained while the gas is being excited via an electrical discharge is briefly discussed in Sec. IIIA.

For the SHRS gain calculations, it was assumed that fluorescence intensities existing in the region of the discharge traversed by the co-propagating laser beams Ω and Ω' are roughly the same as those that would be externally radiated from the surface of a fluorescent lamp having discharge parameters similar to those assumed for the proposed device. In Sec. IIIB this assumption is briefly examined, and is found to be reasonable. It is additionally noted in Sec. IIIB that, via the establishment of EIT, the extent of radiation trapping would be greatly decreased in the region of the discharge traversed by the laser beams. It is likely that this would have the beneficial effect of reducing the electron density in the same region.

In Sec. IIIC, possible experimental configurations with which realization of the device can be attempted are discussed. To be considered satisfactory, any such configuration must include means for injecting an intense, externally or internally generated, "starter" laser pulse into the optical cavity of the dressed-atom laser device. The required intra-cavity-injected energy of the starter pulse would be the energy needed to coherently phase all of the atoms within the vapor cell present in the volume irradiated by the two co-propagating laser beams. In the scientific literature dealing with EIT, this amount of energy is usually termed the "preparation loss".

Section IV briefly summarizes the predicted main characteristics of the laser device proposed in the present paper.

II. NONLINEAR PUMPING OF Λ-TYPE DRESSED-ATOM VAPORS

A. Dressed-atom energy level diagram for Λ atoms

Figure 1 represents the dressed-atom energy level diagram for a gas of coherently phased Λ atoms. Diagrams of this type first appeared and were explained in Ref. [3], an illuminating paper in which it was shown how such diagrams can be simply utilized to gain insight into both coherent and incoherent optical processes occurring in driven three-level systems. Although in Ref. [3] only the cascade-type three-level atomic system is analyzed, extension to the Λ system is trivially simple. It will here be assumed that the reader is familiar with the content of Ref. [3].

From Fig. 1, one easily sees that all atoms in a coherently phased Λ gas must occupy $|2\rangle$ levels, irrespective of the ratio of Rabi frequencies Ω and Ω'. Establishment of EIT in Λ systems is represented in this figure by the indication that an upward transition at either ω_o or ω_o' originating from any $|2\rangle$ level is forbidden. One also sees from this

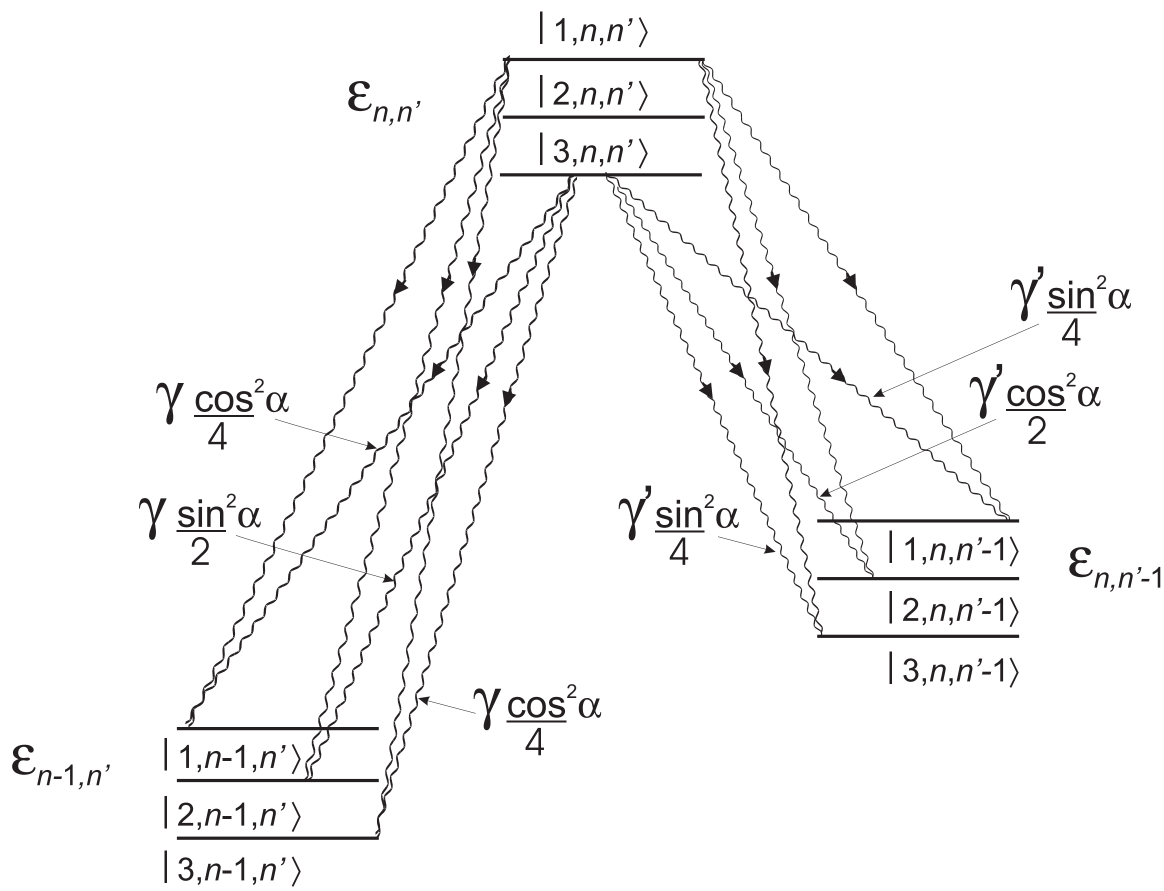

Fig. 1. Diagram showing all spontaneous emission decays which are allowed from the three perturbed states of $\varepsilon_{n,n'}$ to lower multiplicities in Λ-type dressed atoms. Spontaneous emission rates on individual transitions are shown $(\tan\alpha = \Omega'/\Omega)$. The quantities γ and γ' are the radiative decay rates for the two bare-atom transitions.

figure that coherently phased Λ atoms cannot fluoresce – the $|2\rangle$ levels are indeed "dark state" levels, as is well known. The figure also shows that the dressed-atom absorption spectrum consists of a doublet about each of the bare-atom resonance frequencies ω_o and ω_o'. The two components of each absorption doublet are symmetrically disposed in frequency about the corresponding bare atom frequency, with both doublet splittings being equal to the "generalized Rabi frequency" $\Omega_g = \left(\Omega^2 + \Omega'^2\right)^{1/2}$. As the intensities of the beams Ω and Ω' are increased, the doublet splittings increase, but the peak absorption strengths and linewidths of all four doublet components remain largely unchanged, provided that the Rabi frequency ratio Ω/Ω' remains the same. Simple analysis contained in Ref. [3] shows that the relative strength ratio of the two absorption doublets is inversely proportional to the square of the ratio of associated Rabi frequencies. All the above properties of coherently phased three-level systems have been well understood for many years, but the rigorous quantum mechanical analyses [4,5] needed to account for these properties at all power levels of the laser beams that dress the atoms can be quite complex and difficult to follow. Dressed-atom diagrams such as Fig. 1 thus offer a lucid, direct way of gaining insight into the properties of any driven three-level system. However, as explained in Ref. [3], such diagrams rigorously apply only in the "secular limit", *i.e.* when the generalized Rabi frequency Ω_g is much greater than the radiative decay rate of either allowed three-level-atom transition. It will become apparent that, in the operating regime envisioned for the laser device here being considered, this condition is well satisfied.

An immediate conclusion that can be drawn from diagrams such as Fig. 1 is that a CW laser based upon three-level dressed atoms possessing either a cascade-type or V-type energy level structure would be, practically speaking, unrealizable. Unlike coherently phased Λ atoms, coherently phased cascade- or V-type atoms strongly fluoresce. Thus, merely to support a 1-m-long, 1-cm^2-area column of dressed atoms of the latter types would require a continuous application of power in excess of 2MW. Here a dressed-atom density of $10^{15}/\text{cm}^3$ is assumed. By contrast, to support a similar column of dressed Λ atoms would require almost no power, since such atoms do not fluoresce.

B. Linear photonic excitation of dressed Λ atoms

Let us now begin to explore whether one can excite a prepared column of dressed Λ atoms in a manner that results in the addition of photons to the two CW laser beams that dress the atoms – hopefully, in sufficient number to sustain the CW oscillation postulated in the model. Consider first what would happen if one were to apply either coherent or incoherent pumping light resonant with the dressed-atom absorption bands, *i.e.* if one were to attempt to pump the dressed atoms directly via a simple linear excitation mechanism. From Fig. 1, one sees that, in the unit step of such a pumping process, a dressed atom is moved from a $|2\rangle$ level into a higher-lying $|1\rangle$ or $|3\rangle$ level, all levels of the latter type being, of course, unoccupied. However, although any $|1\rangle$ or $|3\rangle$ level is seen to be connected to four other levels of the same type via two allowed optical transitions at ω_o and two at ω_o', a dressed atom excited via linear absorption cannot be driven to any other level by either the Ω or the Ω' laser beam, since the upward and downward transition probabilities exactly cancel. Thus, a dressed atom directly excited via linear absorption cannot provide any optical gain for the Ω and Ω' CW laser beams – it, in fact, produces optical loss for these beams through the following mechanism. The excited dressed atom cannot be driven to any other level by the two laser beams, but it can freely spontaneously emit to lower levels at six different wavelengths. Two of these wavelengths correspond to the excited atom being returned to a $|2\rangle$ level, which would terminate the possibility of any further fluorescence decay events occurring. If all the excited atoms could only spontaneously emit radiation via this channel, the net effect of linear photoexcitation would be simple conversion of pump photons into isotropically emitted spontaneous emission photons. However, as Fig. 1 shows, a dressed atom excited, for example, to a $|1\rangle$ level can also spontaneously emit with equal probability to either a $|1\rangle$ or a $|3\rangle$ level. Since such decay necessarily implies that additional events in the chain of spontaneous emission decays must still occur, and since only the propagating laser beams can supply the pumping energy for any such additional events, one sees that linear photoexcitation of the dressed-atom gas actually produces loss for the propagating laser beams. On the average, for every dressed atom excited via linear photoexcitation, one photon will necessarily have to be removed at some later time (generally, a few radiative lifetimes) from one of the CW laser beams. That it is impossible, in the absence of population inversion, to extract stored energy from the upper atomic level of a driven Λ system using only the self-consistent propagation of the beams Ω and Ω' has been known for a long time. It is clearly stated in Ref. [6], for example.

C. Non-parametric nonlinear photonic excitation of dressed Λ atoms

In view of the analysis given above, one is led to consider pumping the Λ-type dressed-atom vapor of the proposed laser device via a *nonlinear* photoexcitation mechanism. One known such mechanism, which at first glance would seem to have some promise of working in this application, is the canonical (*i.e.* three-photon) form of stimulated hyper-Raman scattering (SHRS) shown in Fig. 2a. Although in all reported experimental studies in which SHRS was realized, one or two co-propagating laser beams provided the excitation, the use of coherent light for pumping this process is not an absolute necessity (Sec. IID). The unit step of the SHRS process shown in Fig. 2a is a simultaneously occurring, energy conserving, three-photon scattering event, in which two nearly resonant pump photons are absorbed, a Stokes-wave photon is emitted, and a ground state atom is excited to the upper level of an allowed two-level transition. Since SHRS is a non-parametric multi-photon process, phase matching is not required. What is most important here is that the light of the two pumping beams be very intense and be spectrally concentrated about the resonances. The fluorescence pump light produced by the electrical discharge in the proposed laser device satisfies both conditions. One thus can consider the possibility that an SHRS mechanism could operate to convert fluorescence photons into photons of the propagating beams at Ω and Ω'. With reference to Fig. 1, a possible unit step of this process, when it is used to excite dressed Λ atoms, might consist of the following events. Fluorescence photons at $\omega_o + \delta$ and $\omega_o - (\Omega_g / 2) - \delta$ would be absorbed, a laser photon at ω_o would be added to the Ω beam, and a dressed atom would be excited from a $|2\rangle$ level to a $|3\rangle$ level, with all events occurring simultaneously. In principle, the SHRS optical gain cross-section, expressed as a function of δ, should be integrated over the entire fluorescence bandwidth (*i.e.* the atomic Doppler width) to account for pumping by all of the fluorescence light. However, from Sec. IIB it follows that there is reason to postpone performing this task. The fact that in the unit step of the SHRS process a $|3\rangle$ level is excited means that the addition of one photon to the Ω beam occurring in this step will at some later time effectively be

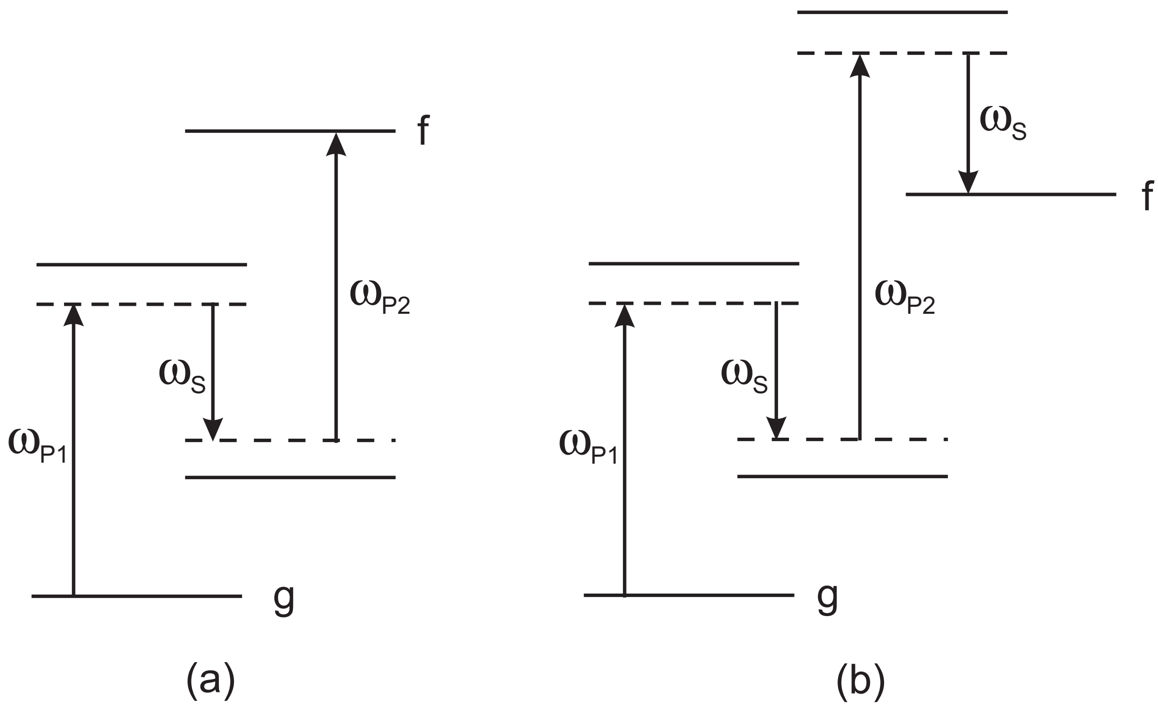

Fig. 2. Schematic diagrams of stimulated hyper-Raman scattering (SHRS) processes: (a) three-photon SHRS, (b) four-photon SHRS.

cancelled by the removal of a photon from either the Ω or the Ω' beam. Thus, nonlinear excitation of a dressed-atom Λ gas via (three-photon) SHRS conversion of highly resonant fluorescence pump light into light of the propagating laser beams provides neither gain nor loss for the latter.

A potential way out of this difficulty is indicated in Fig. 2b, which schematically depicts a *four-photon* SHRS process. The process shown here now becomes of interest for the simple reason that *two* laser photons are produced in the unit step. Thus, even though excitation of a $|1\rangle$ or $|3\rangle$ level again occurs in the unit step, and even though such excitation again implies subsequent loss of a laser beam photon on the average, it is seen that the unit step results in a net gain of one laser beam photon.

Before proceeding further, we should note that a simple energy consideration provides an additional reason why three-photon SHRS could not serve as the pumping mechanism in the proposed device, if the pump light is restricted to be the discharge-produced fluorescence spectrally centered on the transitions at ω_o and ω_o'. In Fig. 3, both the dressed-atom absorption bands and the discharge-produced fluorescence are, for simplicity, represented as having spectral profiles that are rectangular in shape, instead of being Gaussian. In terms of this figure, the limiting condition for operation of the proposed device would evidently be $\Omega_g \geq 2(\Delta\omega_D)$. If Ω_g were less than $2(\Delta\omega_D)$, linear absorption of the pump light by the dressed atoms would occur, resulting in the disastrous effects outlined in Sec. IIB. However, with $\Omega_g \geq 2(\Delta\omega_D)$, it is seen that three-photon SHRS also cannot occur, since in the unit step of such a process, two fluorescence pump photons would be absorbed, one laser photon (at ω_o) would be emitted, and a dressed

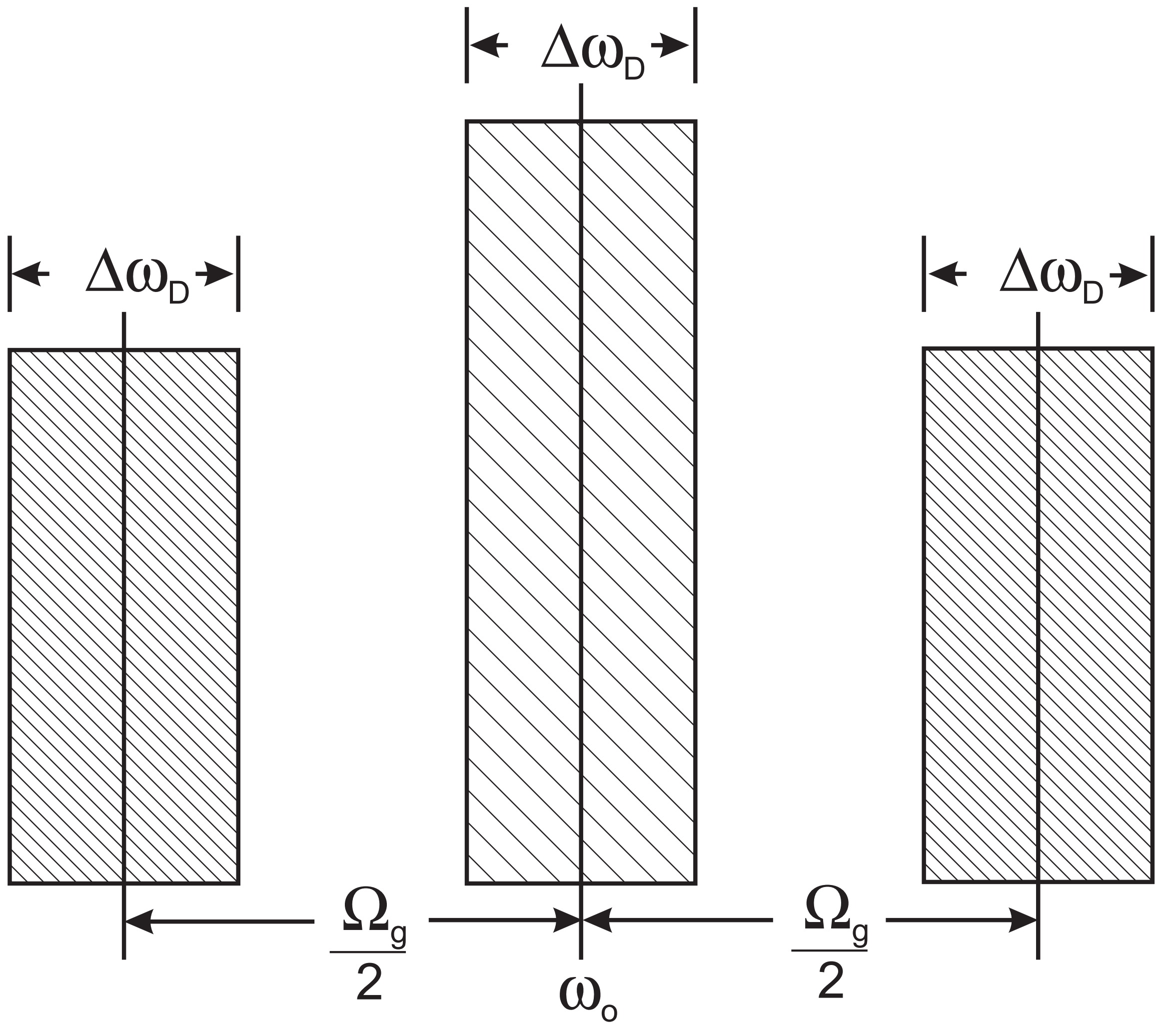

Fig. 3. Simplified representation showing relative spectral positions of dressed-atom absorption bands and discharge-produced fluorescence band occurring at a Λ-type three-level-atom transition.

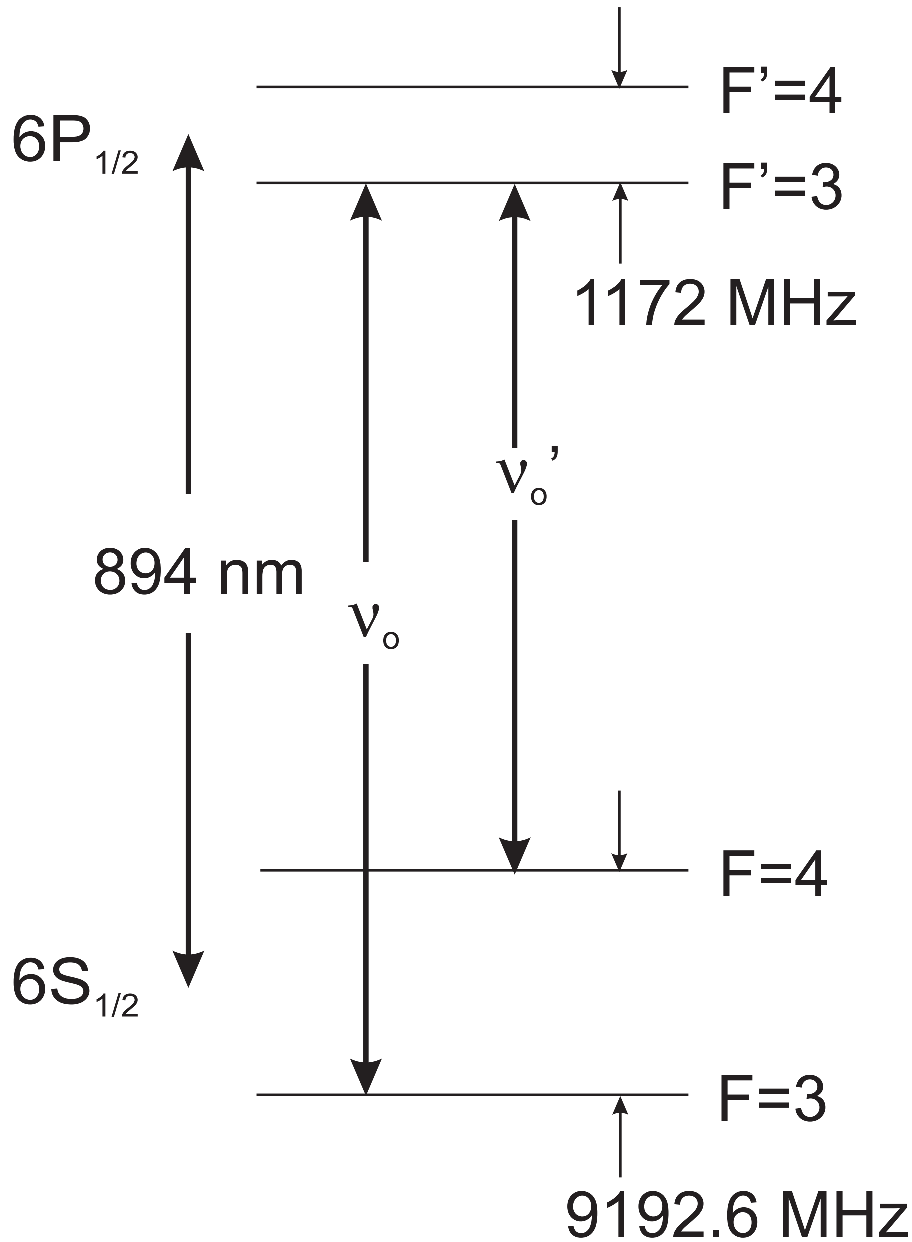

Fig. 4. Lowest energy double-Λ structure in Cs.

atom would be excited from a $|2\rangle$ level to, say, a $|1\rangle$ level higher in energy by $\omega_o + \Omega_g / 2$.

With the energy level structure of dressed atoms, a *four*-photon SHRS process is also possible. Here, in the unit step, two fluorescence pump photons would again be absorbed, *two* laser photons would be emitted, and a dressed atom would be excited from a $|2\rangle$ level of a given multiplicity to, say, the $|1\rangle$ level *of the same multiplicity*. While such a process could in principle provide laser gain, one sees from Fig. 3 that four-photon SHRS could not serve as the pumping mechanism in the proposed device, if, again, fluorescence centered at ω_o and ω'_o were the only available pump light. For dressed-atom excitation via four-photon SHRS to be considered as a potentially feasible pumping mechanism in the proposed device, there must be discharge-produced fluorescence at some other nearby transition(s) that can be used for pumping. Such nearby transitions are provided by the double-Λ structures of alkali atoms. For Cs, the lowest energy double-Λ structure is that shown in Fig. 4. On the basis of this structure, we now proceed to investigate whether enough optical gain can be developed with excitation of Cs dressed atoms via four-photon SHRS to provide sufficient amplification for the two propagating CW laser beams postulated in the model.

D. Optical gain for *n*-photon SHRS excitation of dressed atoms

Figure 5 represents the unit step of a possible four-photon SHRS process by which Cs dressed atoms could be excited in the proposed device. In this figure, it has been arbitrarily assumed that the lasing transitions for the Ω and Ω' beams are those connecting the two $6S_{1/2}$ hyperfine levels with the lower-lying $(i.e.\, F' = 3\,)\, 6P_{1/2}$ hyperfine level. In Fig. 5, the pump light I_P is assumed to be the discharge-produced fluorescence at $\nu_o + 1172$ MHz. It is apparent that, by making these assumptions, one automatically determines the value of the CW laser power circulating within the laser cavity. (Conceptual problems associated with standing-wave nodes are avoided by assuming unidirectional propagation of laser light within the cavity.) The value of Ω_g is here ≈ 2.95 x 10^{10} sec^{-1}. Assuming equal Rabi frequencies $(\Omega = \Omega')$, one has $\Omega \approx$ 2.08 x 10^{10} sec^{-1}. The Cs $6S_{1/2} \leftrightarrow 6P_{1/2}$ transition dipole moment (in SI units, which will be used throughout Sec. IID) is taken to be μ_o= 2.8 x 10^{-29} Cm, the value given in [7]. Since $\mu_o E = \hbar\Omega,$ where E is the peak amplitude of the Ω laser field, one has $E \approx 78{,}000\ \text{Vm}^{-1}$. This corresponds to a cycle-averaged, intra-cavity circulating Ω laser beam power I_S ≈ 8.1 x 10^6 W m^{-2} (810 W cm^{-2}), since $I_S = (1/2)\varepsilon_o c \left|E\right|^2$. The Ω' laser beam would have equal intensity. It is thus seen that, for the proposed device to work, it would necessarily have to operate at a rather high CW power level.

Using standard methods of nonlinear optics, one can estimate the optical gain characterizing an *n*-photon SHRS process. (Ordinary stimulated Raman scattering (SRS) can be viewed as a *two*-photon SHRS process.) The growth of the stimulated Stokes

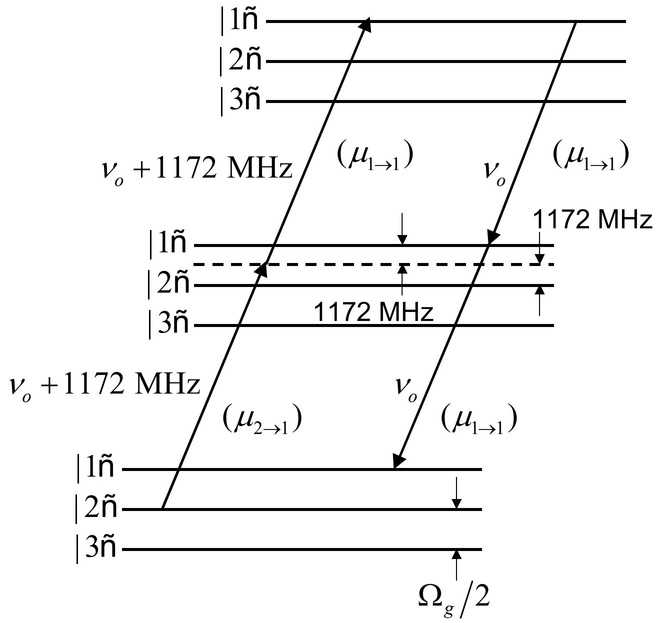

Fig. 5. Unit step of a four-photon SHRS process by which Cs dressed atoms can be excited in the proposed device. Here $\nu_{P1} = \nu_{P2} = \nu_o + 1172$ MHz.

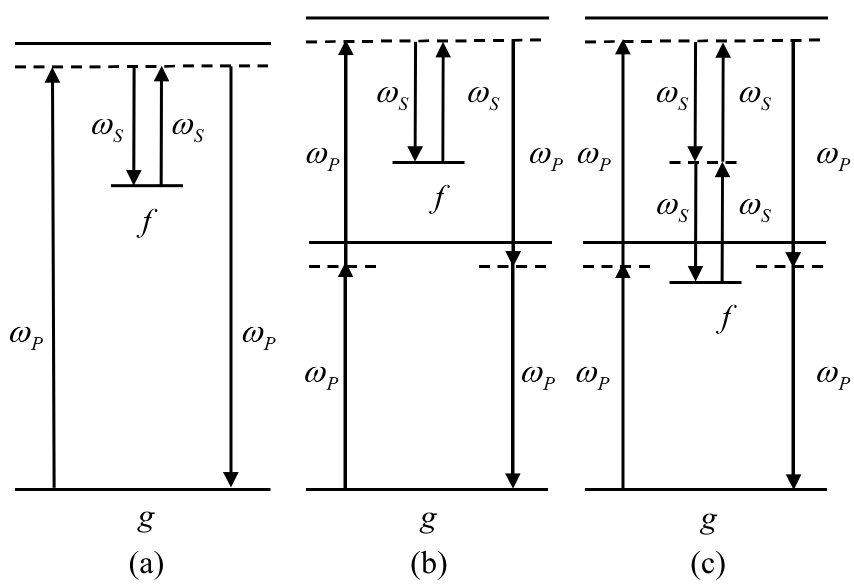

Fig. 6. Schematic diagrams of $2n$-wave photonic processes used for describing Stokes-wave growth in n-photon SHRS processes: (a) 2-photon SHRS (*i.e.* SRS), (b) 3-photon SHRS with a single pump frequency ω_P, (c) 4-photon SHRS with a single pump frequency ω_P.

wave in *n*-photon SHRS can be described entirely in terms of a 2*n*-wave interaction (Fig. 6). As indicated, for example, by the analysis given in Chapter 3 of [8], the growth with distance *z* of the Fourier transform $E_S(z)$ of the electric field of the Stokes wave can be expressed by the equation:

$$\frac{\partial E_S(z)}{\partial z} = \frac{i\omega_S}{2\varepsilon_o c n_S} P_S^{NL}(z)\exp(-ik_S z), \tag{1}$$

where $E_S(z)$ is the slowly varying part of the total Fourier transform $E_S(z)\exp(ik_S z)$. The quantity $P_S^{NL}(z)$ is the total Fourier transform of the nonlinear polarization at ω_S, and is given by:

$$P_S^{NL}(z) = \varepsilon_o K \chi^{(n)}(-\omega_S; \omega_1, \omega_2, ...\omega_n) E_1^{(*)}(z) e^{\pm ik_1 z} E_2^{(*)}(z) e^{\pm ik_2 z} ... E_n^{(*)}(z) e^{\pm ik_n z}. \tag{2}$$

In this equation, the value of *K* depends upon the order of the nonlinear process and on the number of distinguishable permutations of the $\omega_1 ... \omega_n$. (A frequency and its negative are considered distinguishable.) Values of *K* for various specific nonlinear processes are listed in Table 2.1 of [8]. For SRS (Fig. 6a), *K*=3/2. For 3-photon SHRS with a single pump frequency ω_P (Fig. 6b), *K*=15/8. For 3-photon SHRS with two distinct pump frequencies ω_{P1} and ω_{P2}, *K*=15/2. For 4-photon SHRS with a single pump frequency ω_P (Fig. 6c), *K*=$(7!/2x2x2)/2^6 \approx 10$.

In Eq. (2), the superscripts $(*)$ on the slowly varying parts of the Fourier transforms imply that the complex conjugate is to be employed if the wave appears in emission in diagrams such as those in Fig. 6. When the superscript $(*)$ applies, one must also use the (-) sign in $\exp(\pm ik_n z)$. In Eqs. (1) and (2), propagation of all waves along the positive z-axis is assumed.

The complex nonlinear susceptibility $\chi^{(n)}$ in Eq. (2) is given by the expression:

$$\chi^{(n)}(-\omega_S; \omega_1, \omega_2, ...\omega_n)$$

$$= \frac{N}{n!\hbar^n \varepsilon_o} \Im_T \sum_{gb_1...b_n} \frac{\mu_{gb_1}\mu_{b_1b_2}...\mu_{b_ng}}{(\Omega_{b_1g} - \omega_1 ... - \omega_n \pm i\Gamma_{b_1})(\Omega_{b_2g} - \omega_2 ... - \omega_n \pm i\Gamma_{b_2})...(\Omega_{b_ng} - \omega_n \pm i\Gamma_{b_n})}, \tag{3}$$

where the summations $\sum_{gb_1...b_n}$ are over all the atomic states. The quantity μ_{ij} is the atomic dipole moment for the transition between states *i* and *j*. In Eq. (3), $\Im_T$ is a permutation operator requiring that the expression following it be summed over all permutations $-\omega_S, \omega_1...\omega_n$. In principle, there are therefore $(n+1)!$ expanded terms in Eq. (3), although many of the terms will be the same if some of the frequencies are equal.

In Eq. (3), the $\Gamma's$ are damping terms introduced to prevent infinite singularities in $\chi^{(n)}$ from occurring when exact resonances are present, that is, when some of the factors in the denominator of Eq. (3) would otherwise be identically zero. Nonlinear optics provides a simple "rule" (*c.f.* p. 35 of Ref. [8]) for determining the choice of sign appearing before each $i\Gamma$ in Eq. (3). This rule states that if the resonant denominator appears to the right of the $\omega_\sigma (\equiv \omega_S)$ arrow, take $-i\Gamma$; if it is to the left, take $+i\Gamma$.

Bearing in mind the above remarks, one can proceed to roughly estimate the nonlinear Λ dressed-atom susceptibility $\chi^{(7)}(-\omega_s;\omega_p,\omega_p,-\omega_p,-\omega_p,\omega_s,-\omega_s,\omega_s)$ that should apply in the present case. It is here simplest to consider initially just the contribution to $\chi^{(7)}$ arising from one specific ordering of the eight waves participating in the unit step of the four-photon SHRS mixing process. This is the ordering shown in Fig. 6c, which corresponds to the dressed-atom four-photon SHRS excitation process shown in Fig. 5. Since the permutation operator $\Im_T$ in Eq. (3) would generate 2x2x2x2 = 16 of these terms, one should enter this number as a multiplying factor in the right hand side of Eq. (1). Entering also the four-photon SHRS *K* factor discussed earlier, and the factor (1/2) already appearing in Eq. (1), one then finds (1/64) to be the purely numerical factor one should use in Eq. (1) in the present case.

In evaluating the sum $\sum_{gb_1..b_n}$ appearing in Eq. (3) for the specific nonlinear excitation scheme shown in Fig. 5, one can note in passing an important new feature that arises because dressed atoms are involved. From this figure, one might initially suspect that, due to quantum interference, the transition probability for the four-photon process depicted might be somewhat reduced by virtue of the fact that the effective intermediate state for the first transition shown could be either a $|1\rangle$ or a $|3\rangle$ level, with the former lying above, and the latter lying below, the virtual intermediate state for the same transition. If the dipole moments $\mu_{2\to1}$ and $\mu_{2\to3}$ had the same sign, this would lead to partial cancellation of the overall four-step transition probability. However, from Ref. [3] one sees that $\mu_{2\to1} = -\mu_o \sin\alpha/\sqrt{2}, \mu_{2\to3} = \mu_o \sin\alpha/\sqrt{2}$ ($\tan\alpha = \Omega/\Omega'$), so that the transition probability contributions arising from the $|1\rangle$ and $|3\rangle$ intermediate state levels do not cancel one another, but in fact additively combine. The same remark would also hold in the case of the second transition. Here, either the $|1\rangle$ or $|3\rangle$ levels of the next highest multiplicity could again, technically speaking, serve as intermediate states, although in this case the virtual intermediate state is resonant with the $|1\rangle$ level and is offset by Ω_g from the $|3\rangle$ level, thereby greatly reducing the relative contribution to the transition probability by the latter. At any rate, the corresponding transition probability contributions would again additively combine, since $\mu_{1\to 1} = \mu_o \cos\alpha/2, \mu_{1\to 3} = -\mu_o \cos\alpha/2.$ In brief, no transition probability cancellation due to quantum interference arises when *n*-photon SHRS excitation is applied to a dressed-atom gas. This example illustrates one of the many important advantages gained in utilizing nonlinear excitation of dressed-atom vapors to generate laser beams at different frequencies. That such advantages automatically accrue with use of this generation technique was first theoretically predicted in Ref. [9]. As illustrated by the example discussed in the Introduction, many subsequent experimental studies have demonstrated that high efficiencies of laser beam generation can indeed be attained when this general scheme is employed.

Assignment of denominator factors in Eq. (3) for the transitions shown in Fig. 6c is straightforward. For the first and seventh transitions, the offset is real (2π x 1172 MHz = 7.36 x 10^9 sec^{-1}), and one can realistically take $\Gamma = 0$. There is an exact resonance at both the second and sixth transitions. By virtue of the "rule" mentioned above, the denominator factor corresponding to the first of these is $(+i\Delta\omega_D)$, while that corresponding to the second is $(-i\Delta\omega_D)$, with $\Delta\omega_D$ being 2π times the Cs-atom Doppler width $\Delta\nu_D$. Similar denominator factors apply in the cases of the third and fifth transitions. The denominator factor associated with the resonance at the fourth transition would be $(+i\Delta\omega_D)$. From Eq. (3), one sees that the combined effect of all these denominator factors is to make $\chi^{(7)}$ a *negative imaginary* quantity. Therefore, with exactly the same reasoning used in nonlinear optics to show optical gain is present in two- and three-photon SHRS, one sees from Eqs. (1) and (2) that it also should be present in four-photon SHRS excitation of dressed atoms. For this process, the power gain is finally seen to be:

$$G_{HR} = \frac{N\omega_S I_P^2 I_S \mu_{2\to 1}^2 \mu_{1\to 1}^6}{8\varepsilon_o^4 c^4 \hbar^7 (\Delta\omega_D)^5 (7.36x10^9)^2}. \qquad (4)$$

The hyper-Raman gain given by Eq. (4) represents an exponential intensity gain per unit length, that is, in the absence of pump power depletion and/or saturation in the efficiency of the SHRS process, the intensity of the laser beam at ω_S would increase by a factor $e^{G_{HR}l}$ in traveling a distance l. The quantity I_P appearing in this equation is the intensity of isotropically radiated fluorescence pump light at ν_o+1172 MHz present everywhere in the region of the vapor cell where the laser beams propagate. We here take this quantity to be 0.1W cm^{-2} (1000 W m^{-2}). For the approximate temperatures (T$\approx$500K) at which it is envisioned that a Cs dressed-atom laser would operate, $\Delta\nu_D \approx 500$ MHz. Therefore, $\Delta\omega_D \approx 3.14$ x 10^9 sec^{-1}.

Let it be initially assumed that the Cs atomic density in the vapor cell is 10^{21} m^{-3}. Assuming laser operation with $\Omega = \Omega'$, the absolute values of $\mu_{2\to 1}$ and $\mu_{1\to 1}$ would then be $\mu_o/2$ and $\mu_o/(2\sqrt{2})$, respectively. The value of μ_o was noted earlier. The angular frequency ω_S $(\equiv \omega_o)$ is here 2.2 x 10^{15} $\sec^{-1}$. Substituting all the various above parameter values, one finds the power gain coefficient for four-photon SHRS calculated via Eq. (4) to be $G_{HR} \approx 0.004$ m^{-1}. Accounting for the effect of fluorescence pump light at $\nu_o' +$ 1172 MHz driving a photonic scheme similar to the one shown in Fig. 5, with both photons still being emitted at ν_o, doubles the value of G_{HR}, making it ≈ 0.008 m^{-1}. In addition, there would be a contribution to the gain at ν_o arising from a four-photon SHRS scheme analogous to the one in Fig. 5, with pumping still occurring at $\nu_o + 1172$ MHz, but with photons at both ν_o and ν_o' now being emitted. Finally, there would be a contribution from a scheme similar to this last one, but with pumping occurring at $\nu_o' + 1172$ MHz. Each of these latter schemes would contribute 0.002 m^{-1} to the gain, making the final estimated total power gain $G_{HR} \approx$ 0.012 m^{-1}. This would allow CW oscillation to occur in a 1-m-long cavity equipped with end mirrors not transmitting a total of more than 1%. If the Cs-atom density were raised to 10^{16} cm^{-3}, the calculated gain coefficient would be much larger (≈ 0.1 m^{-1}). On the other hand, having ten times more Cs atoms would imply a ten times larger "preparation loss" (Sec. IIIC), making it more difficult to "start" the dressed-atom laser.

While the results of the above gain estimates would seem to imply that the proposed laser device should be feasible, one should bear in mind that many approximations were involved in making in these estimates. For example, the fluorescence pump power I_P at each of the double-Λ transitions in Fig. 4 was assumed to be 0.1 W/cm^2. Had this quantity been taken to be significantly different, the gain G_{HR} would have been much different, since $G_{HR} : I_P^2$. For reasons such as this, the calculated optical gain for the device should probably be viewed as being only *marginally sufficient* for CW laser action to occur.

A more serious question regarding the feasibility of the device is associated with the assumption made above that a four-photon SHRS dressed-atom excitation process could be driven by incoherent fluorescence pump light. While it is beyond the scope of the present discussion to present a complete argument that this is possible, we here note that there is compelling evidence in the scientific literature that pumping of SRS (*i.e. two*-photon SHRS) can, under certain conditions, be done as effectively with broadband incoherent light as with narrow-band laser light. For example, the results of a theoretical and experimental investigation of SRS in the field of a Gaussian noise pump in a medium with active length l are presented in Ref. [10], where it is shown that the nature of the stimulated scattering depends upon the relationship of the correlation time τ_{cor} of the noise-pump field, on the one hand, and the transverse relaxation time T_2, and the characteristic group-lag time T_3, on the other hand. For co-propagating pump and Stokes waves, $T_3 = l\left(\frac{1}{u_P} - \frac{1}{u_S}\right)$, where u_P and u_S are the respective group velocities. For opposing waves, $T_3 = l\left(\frac{1}{u_P} + \frac{1}{u_S}\right) \approx \frac{2l}{c}$. In atomic gases, SRS typically occurs in the forward direction when the regime $T_2 > \tau_{cor} > T_3$ is satisfied. It is shown theoretically in

[10] that, for this regime, a broadband noise pump is as effective as a coherent monochromatic pump having the same intensity. The same result would seem also to apply when dressed atoms are excited via four-photon SHRS, as in the proposed laser device, the reason being that T_3 would here be almost identically zero, since pump and Stokes waves are co-propagating. In the proposed device, all pump and Stokes-wave frequencies occur within the two narrow spectral regions in which EIT is established, *i.e.* the spectral intervals $|\omega - \omega_o| \le (\Omega_g - \Delta\omega_D)/2$ and $|\omega - \omega_o'| \le (\Omega_g - \Delta\omega_D)/2$. As can be seen from plots of χ', the real part of the dressed-atom linear susceptibility (*c.f.* Fig. 2 of [9], or Fig. 13 of [11]), the slope of this function changes very little in the above spectral intervals, implying constant group velocities over the same intervals.

In the theory presented in Ref. [10], it is deduced that, in the regime $T_2 > \tau_{cor} > T_3$, the linewidth of the SRS Stokes wave generated is approximately that of the stochastic pump light. This would imply that the spectral width of laser beams generated in the proposed Cs discharge-pumped, dressed-atom laser device should always be approximately 500 MHz.

There is a further interesting possibility connected with the dressed-atom excitation scheme that has here been presented. Consider together both Figs. 4 and 5. It has been shown that lasing could possibly occur on the two lower-leg transitions of the double-Λ structure when pumped by fluorescence at the two upper-leg transitions. In the same way, one should expect that lasing could occur on the two upper-leg transitions, with the pump power being the fluorescence at the lower-leg transitions. Therefore, in principle, lasing could also *simultaneously* occur on all four double-Λ transitions, with one generalized Rabi frequency, Ω_g = 2.95 x 10^{10} sec^{-1}, characterizing lasing on both upper- and lower-leg transitions. In [12] and [13], it is shown theoretically that propagation of pairs of optical pulses in a double-Λ configuration is especially favorable for the attainment of EIT at all four frequencies. The high conversion efficiency attained in the mixing experiment of [2], mentioned in the Introduction, may partly have resulted from this effect.

For the sake of completeness, dressed-atom excitation via parametric nonlinear processes will now briefly be considered.

E. Parametric nonlinear photoexcitation of dressed Λ atoms

The four-wave mixing scheme shown in Fig. 7 represents a particularly simple and symmetrical example of such a process. Here it is assumed that, in the unit step, pump photons at $\omega_1 = \omega_o + \delta$ and $\omega_2 = \omega_o - \delta$ are absorbed, while two photons at ω_o are simultaneously emitted. The contribution to the four-photon-process transition probability by the two narrow-band light sources at $\omega_o \pm \delta$ must be integrated over the entire frequency range $0 \le \delta \le \Delta\omega_D/2$ to account for pumping by all of the fluorescent light. Assuming that all four "waves" are co-propagating, the process is automatically perfectly phase matched for any value of δ. Via EIT, there is also ideally perfect transparency at ω_o. For the same reason that was shown earlier to apply when dressed

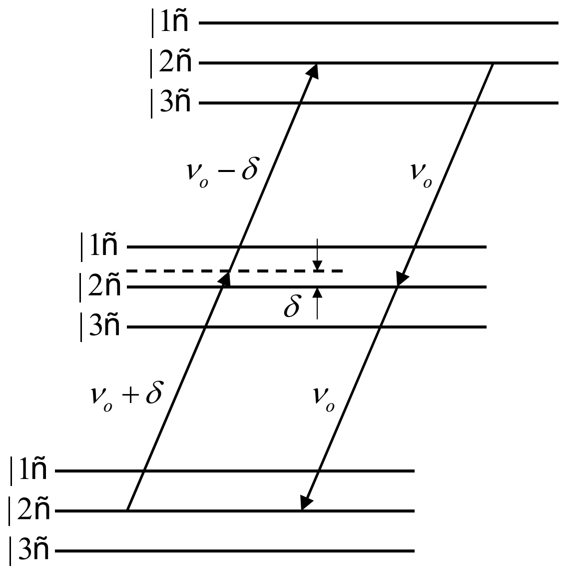

Fig. 7. A possible four-wave parametric process for nonlinearly exciting dressed Λ atoms.

atoms are excited via SHRS, in Fig. 7 the four-photon unit step transition probability is not reduced by quantum interference, *i.e.* at each of the four photonic transitions making up the unit step process, the $|1\rangle$ and $|3\rangle$ intermediate state levels both contribute equally to the overall transition probability, even though they are spectrally located equidistantly on either side of the corresponding virtual intermediate state. In addition, when Λ dressed atoms are excited by the four-wave mixing process in Fig. 7, the atoms never leave their starting $|2\rangle$ levels. Therefore, no requisite subsequent loss of a laser photon has to accompany each unit step event, unlike what necessarily has to occur in all orders of multi-photon SHRS. In fact, two laser photons are produced in the unit step process of Fig. 7. Finally, the parametric process of Fig. 7 can occur for any value of Ω_g, unlike what was seen to hold for n-photon SHRS, where, for any order n, there exists a value of Ω_g above which the process cannot physically occur, when the pump light is restricted to occur around ω_o and ω_o'. In brief, many advantages would appear to accrue with use of four-wave mixing processes to excite dressed atoms in the proposed laser device. Therefore, let us now attempt to probe somewhat further this potential.

The problem with this scheme immediately becomes apparent when one combines Eqs. (1) and (2) and applies the result to this system. The following equation results:

$$\frac{\partial E_S(z)}{\partial z} = \frac{i3\omega_S}{4cn_S}\chi^{(3)}(-\omega_S;\omega_1,\omega_2,-\omega_S)E_1E_2E_S^*(z) = \beta E_S^*(z). \tag{5}$$

For $E_S(z)$, the slowly varying part of the Fourier transform of the Stokes field, to experience gain, the quantity β must be a *fixed* positive real number. However, β is proportional to the product of the slowly varying parts of the Fourier transforms of the noise-pump fields at ω_1 and ω_2, and there is no way that it can be regarded as being a constant function of z. This is simply a manifestation of the well known fact that in parametric processes, phase relationships with the driving pump sources are important, while in non-parametric processes, only pump intensities matter.

III. EXPERIMENTAL APPROACHES FOR TESTING FEASIBILITY OF THE PROPOSED LASER DEVICE

A. Containment vessel for the atomic vapor and electrical discharge

To maintain a stable, arbitrarily long, isothermal column of pure Cs vapor, at pressures adjustable anywhere over a wide range (: 0.01- 100 Torr), and to provide a means for producing throughout most of this column intense $6S_{1/2}-6P_{1/2}$ fluorescence via electronic impact excitation, one can effectively utilize what is termed a *heat-pipe discharge tube (HPDT).* Such robust devices were first proposed [14] and experimentally demonstrated [15] more than thirty years ago. A schematic diagram of an HPDT is shown in Fig. 8. The principles by which such a device operates when the continuous electrical discharge occurs axially are discussed in [15]. If the necessity should arise in the laser device under consideration to keep the gas volume in which the discharge occurs separate from the one in which the Ω and Ω' beams propagate, one could perhaps utilize an HPDT with a transversely occurring discharge (Fig. 9) – although the authors are unaware of reports where such an HPDT was actually constructed. Throughout Sec. IIIB, an axially occurring discharge will be assumed.

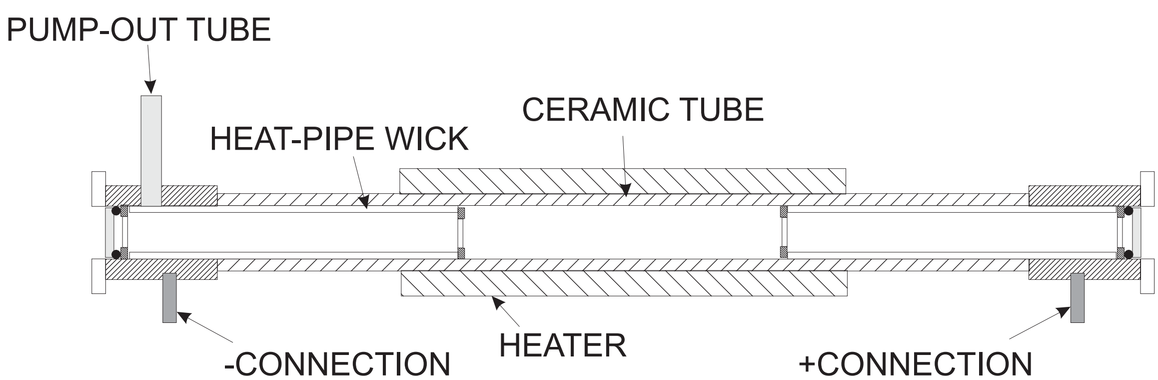

Fig. 8. Schematic diagram of a heat-pipe discharge tube (HPDT) with an axially-occurring discharge.

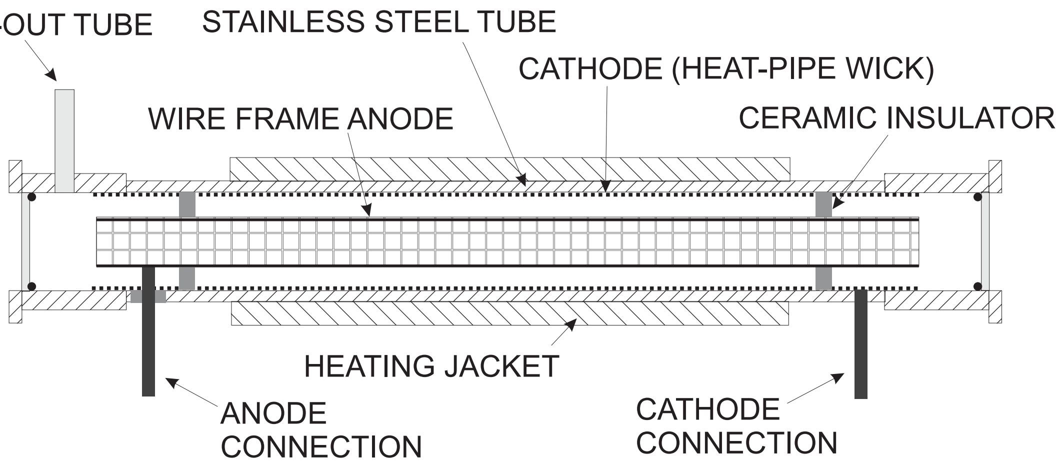

Fig. 9. Schematic diagram of an HPDT with a transversely-occurring discharge.

B. Radiation trapping in atomic vapor electrical discharges

Theories of radiation trapping occurring in low-pressure gaseous electrical discharges (see, for example, Ref. [16]) generally assume that the following canonical event continually occurs. A ground state atom in the discharge absorbs a resonance photon and becomes excited to the upper state, remains there for some time, and eventually releases the photon again. This absorption/reemission process can be repeated many times, until the photon finally escapes from the vapor. The externally radiated CW power of the discharge is proportional to the product of the average excited state density and the average rate at which an excited atom in the discharge creates photons that are incident on the inner surfaces of the gas-containing vessel (*i.e.* that are leaving the vapor). In the case of no trapping, this second factor would simply be γ, the spontaneous emission decay rate of the atom. However, if each photon is absorbed and reemitted g_o times, the effective rate at which photons escaping from the vapor are produced by each excited atom becomes γ / g_o. For the CW light output of the discharge to be roughly the same as when no trapping occurs, the average excited state density must therefore here exceed the value in the no-trapping case by a factor g_o. In the scientific literature dealing with low pressure gas-discharge lamps, g_o is usually termed the *trapping factor.* It can often have a value as high as : 10^4 [16]. Thus, according to prevalent current models of radiation trapping in gas discharges, most of the energy produced via electron impact with atoms of the gas resides in the medium as atomic excitation, with purely photonic (*i.e.* electromagnetic) energy representing only a tiny fraction of the stored energy density. However, irrespective of the amount of trapping, the isotropically propagating fluorescence photon flux inside a gas discharge would always have a value roughly the same as that which would be externally radiated from the surface of a fluorescent lamp constructed from such a discharge. This was implicitly assumed in the calculations of optical gain made in Sec. IID.

In principle, a different model from the one discussed above can also describe low pressure gaseous discharges. As discussed in various texts (*e.g.* [17]), when there is no collision broadening, the Lorentzian cross-section for scattering by a two-level atom exclusively applies to the elastic (*i.e.* Rayleigh) component. The probability of inelastic scattering is zero, *i.e.* no real linear excitation of the atom occurs. In such a regime, photons produced via electron impact would not be absorbed by the atoms of the gas, but would simply propagate throughout the latter in a random walk manner. If this elastic regime were to prevail in the proposed laser device, significant enhancement of the fluorescent photon density in the region traversed by the laser beams could in principle occur (*vide infra*), which in turn would lead to much higher values for G_{HR} than those calculated in Sec. IID. However, since the cross-sections for resonance broadening collisions are extremely large, one can easily see why it is generally believed that scattering of fluorescence photons by the atoms in a low pressure gas discharge must occur inelastically. For example, the resonance broadening cross-section for atomic thallium involving transitions from the ground state to the $7S_{1/2}$ state is σ_{Tl-Tl} = 5.38 x $10^{-8}/\overline{v}$ cm^2 [18], where $\overline{v}$ is the relative velocity of the colliding particles. If the same cross-section applied to the Cs $6S_{1/2} - 6P_{1/2}$ case, this would imply a Cs-atom collisional rate : 0.5 x 10^8 sec^{-1} at N_{Cs} = 10^{15} cm^{-3}. Thus, with nearly unity probability, an excited Cs atom would make a collision with another atom of the gas before it radiatively decays.

On the basis of either the inelastic or elastic scattering models for radiation trapping in discharges, the following interesting effect could possibly occur in the proposed device. Because EIT would be established throughout the region of the vapor cell where the Ω and Ω' beams propagate, no radiation trapping of fluorescent light would be expected to occur in this region. There would in principle here be no Cs excited state population, and, as a consequence, no extra ionization could be produced via electron impact excitation of excited state atoms. There is some experimental evidence (*e.g.* Fig. 18.5 of [16]) showing that radiation trapping can have a strong effect on the electron density in gas discharges. The trapping increases the electron temperature, and tends to make the electron energy distribution more similar to the equilibrium value (*i.e.* the Saha distribution). When an axial discharge is employed in the proposed laser device, this effect might tend to reduce the electron density in the region where the laser beams actually propagate. This would tend to lessen any disruptive effect electrons in the discharge might have on the purely photonic processes involved in dressed-atom lasing.

We conclude Sec. IIIB by describing a scenario through which the steady-state fluorescence photon density could become enhanced in the region of the vapor cell (hereafter termed the "core") where the laser beams propagate in the proposed device. To keep the discussion simple at the outset, it will here be assumed that the elastic scattering model outlined above applies. Consider photons, created via electron impact excitation outside the core, diffusing inwardly towards the latter. At the cylindrical boundary surface, these photons readily enter the core, and then freely propagate, since elastic scattering no longer occurs in this region. (Elastic scattering would occur at the dressed-atom absorption band frequencies, but these are shifted away from the atomic resonance lines by $\pm\Omega_g/2$. The fluorescence photons entering the core are assumed to have a Doppler spectral distribution centered about the atomic resonance lines.) On the other hand, when fluorescence photons inside the core strike the cylindrical boundary surface, they are largely elastically scattered back into the core. This "one way" photon transmittance at the core boundary could thus potentially lead to very large fluorescence photon density enhancements within the core.

Assuming more realistically that only inelastic scattering occurs in the above scenario, one can still argue that the steady-state fluorescence photon density would become enhanced in the core. However, the authors have not performed calculations to confirm that this effect should occur.

C. Cavity configurations for "starting" discharge-pumped dressed-atom lasers

Thus far in the present paper, it has been assumed that a pre-existing column of Cs dressed atoms is nonlinearly driven by fluorescence pump light, resulting in the continual generation of enough additional laser photons to replace those which leave the laser cavity in the form of the CW laser output beams produced by the device. It was explained in Sec. IIA that, once coherently phased, a gas of Λ-type atoms essentially does not require additional power to be expended in order for it to remain in this prepared state. However, to initially form, for example, a 1-m-long column of coherently phased Cs atoms at a density of 10^{15} cm^{-3} would require a total "preparation energy" [19] of roughly 20 mJ cm^{-2}. This energy would have to be supplied to the Cs vapor cell in the device through direct irradiation of the latter by simultaneous pulsed laser beams at ω_o and ω_o', with the energy of each beam being at least 10 mJ cm^{-2}. Since this value represents an intra-cavity-injected energy, a somewhat challenging technical problem presents itself here, in view of the conclusion made in Sec. IID that the total transmittance of both cavity mirrors should not exceed : 1%. A rather extensive amount of equipment would apparently be needed to "start" a Cs discharge-pumped laser in this manner. This equipment is diagrammed in Fig. 10. In initial attempts to demonstrate operation of the device, it might be best to work with small area beams, e.g. : 0.1 cm^2. The total starting pulse energy applied to the 1% transmitting mirror in Fig. 10 would then have to be : 200 mJ. Although the frequencies of the two starting laser beams should be accurately centered at ω_o and ω_o', no purpose would be served in making the line widths of these beams less than about 500 MHz. Although there would exist some latitude in the choice of starting laser pulse widths, a measure of guidance is provided by considerations such as those that were presented in Sec. IID. For example, it would probably be best for the circulating intra-cavity power of each of the pulsed starting laser beams to exceed the 810 W/cm^2 CW steady-state value expected in the case of a Cs dressed-atom laser. In a 0.1 cm^2 beam, therefore, each starting laser beam power should be no less than : 100 W. This would imply a starting laser pulse width : 10 μsec. After making a final pass through the Ti-sapphire amplifier rod in Fig. 10, each starting laser beam should therefore have a power : 10kW. This power level should be achievable by multi-passing the modulated CW Ti-sapphire laser beam four or five times through the Ti-sapphire amplifier rod, assuming that the latter is pumped by a relatively fast (: 10 μsec) flashlamp. In this way, in principle, Cs CW dressed-atom laser action could be started on either of the Cs double-Λ transition pairs.

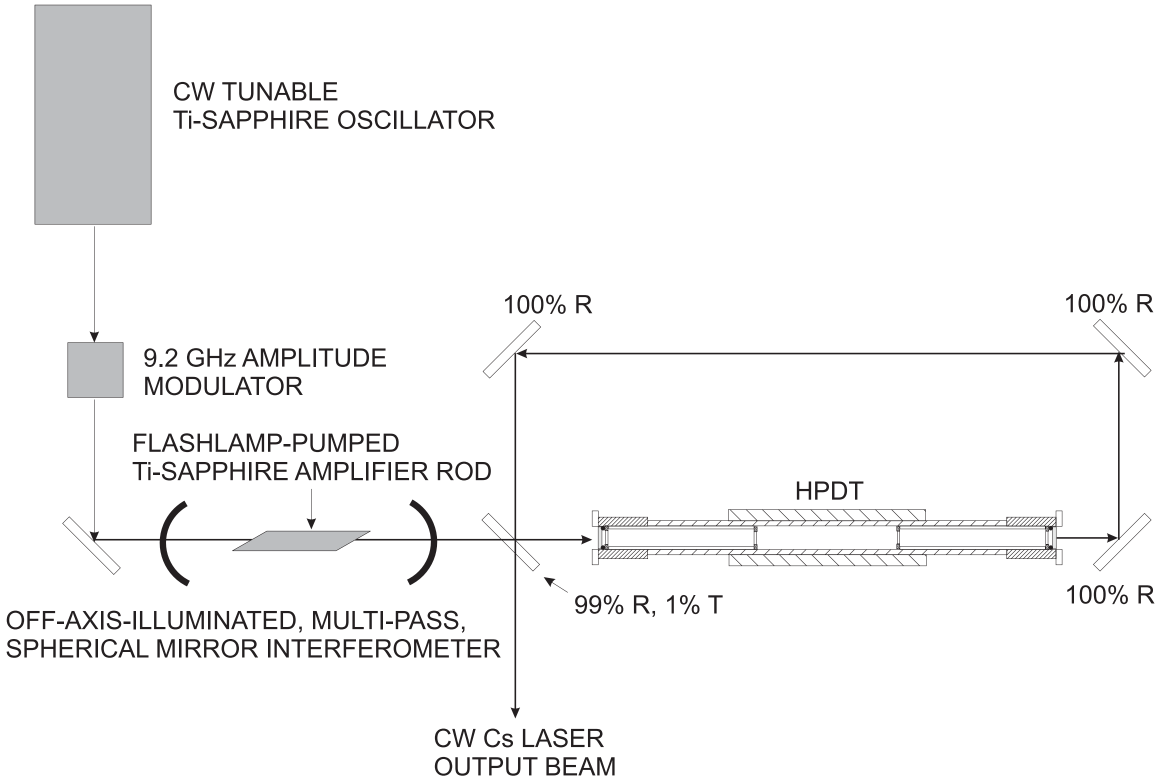

Fig. 10. Diagram of apparatus that can be used to initiate CW operation of a Cs discharge-pumped dressed-atom laser.

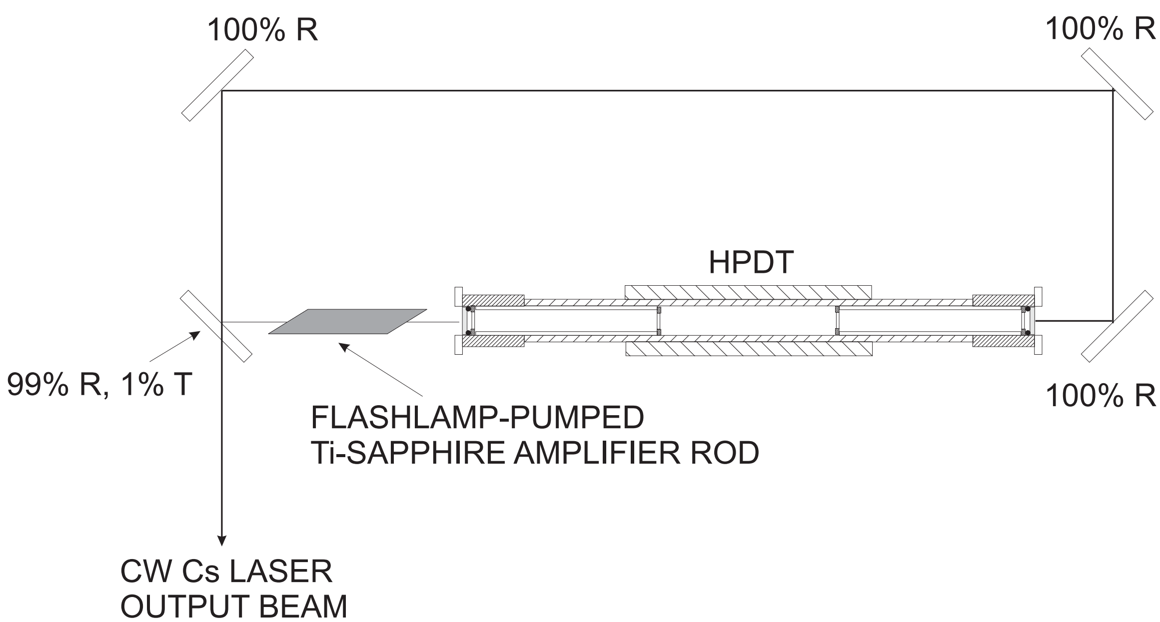

Fig. 11. A possible alternative configuration for initiating CW operation of a Cs discharge-pumped dressed-atom laser.

By tuning the CW Ti-sapphire laser so that the amplitude modulation sidebands become resonant with the other double-Λ transition pair, and once more pulse pumping the Ti-sapphire amplifier, one should in this way be able to induce CW laser action to occur on all four Cs $6S_{1/2} - 6P_{1/2}$ optical transitions.

If starting the proposed dressed-atom laser were to always require the use of equipment as complex as that shown in Fig. 10, the laser, even if it were realizable, would clearly represent a rather impractical device. The authors believe that the proposed laser could be started in a much simpler way, using a configuration similar to the one sketched in Fig. 11. Here the Ti-sapphire rod is simply incorporated into the Cs-laser cavity as an extra element. No basis for tuning is provided, in order to avoid any significant additional optical cavity loss. When the rod is pumped by an external light pulse, very strong Ti-sapphire laser action will occur over a very wide band of frequencies that would normally include all four Cs $6S_{1/2} - 6P_{1/2}$ optical transitions. In a gas of randomly phased atoms, these transitions would of course all be strongly absorbing, and this would initially prevent the pulsed Ti-sapphire lasing from occurring exactly at the resonance frequencies. However, strong Ti-sapphire lasing would occur in spectral regions that closely abut the Cs resonances. Via various n-photon SHRS processes driven by the intense Ti-sapphire laser light, nonlinear excitation of the Cs atoms conceivably could then happen, leading to the eventual establishment of discharge-pumped, CW dressed-atom laser action on all four Cs transitions as the Ti-sapphire laser pulse energy circulating in the cavity slowly decays. An advantage of this scheme is that it is relatively simple and inexpensive.

IV. SUMMARY

In the present paper, the feasibility of constructing a discharge-pumped, CW dressed-atom gas laser has been examined. It is concluded that such a device could be made to operate when a meter-or-so-long column of coherently phased Cs atoms becomes additionally excited by the nonlinear photonic process of four-photon stimulated hyper-Raman scattering (SHRS). Via this process, discharge-produced fluorescence at all four allowed $6S_{1/2} - 6P_{1/2}$ optical transitions can become efficiently converted to CW laser radiation at the same four frequencies. Through this pumping scheme, the value of the intra-cavity circulating CW beam power at each of the four frequencies is determined to be : 810 W/cm^2. Since estimates indicate that, in a single pass through the vapor cell, each laser beam can potentially be amplified by about 1%, this suggests that a total CW laser output power of up to : 32 W/cm^2 could be extracted from the device. To initiate CW lasing of the device on either the Cs upper- or lower-leg transition pair would require simultaneous injection into the laser cavity of two : 10-μsec-long resonantly tuned laser pulses, each with an intra-cavity energy of at least 10 mJ/cm^2.

Finally, the authors would like to comment briefly on their choice of the word "phasers" appearing in the title of the present paper. For more than a dozen years, M.O. Scully has suggested (*e.g.* in [11]) that this word appropriately describes laser devices that operate on the basis of coherently phased three-level atoms. We feel that this suggestion of Scully's makes good sense. In addition, since the device here being proposed operates by directly converting discharge-produced fluorescence (*i.e.*

incoherent or "unphased" light) into laser light (*i.e.* coherent or "phased" light), the use of the word "phasers" here seems doubly appropriate.

ACKNOWLEGMENTS

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